

Searching for axion and dark photon dark matter using dielectric haloscopes

Dissertation zur Erlangung des Doktorgrades an der Fakultät für
Mathematik, Informatik und Naturwissenschaften
Fachbereich Physik der Universität Hamburg

vorgelegt von
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Hamburg
2026

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05.06.2026

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Abstract

The axion and dark photon are two popular dark matter candidates with expected masses below 1 eV. The axion was originally predicted by a solution to the strong CP problem, that is, the unexpected absence of charge-parity (CP) violating effects in quantum chromodynamics (QCD). Its small coupling to the Standard Model and simple production mechanism sparked a strong interest in the axion as a dark matter candidate and various experiments started searching for it. An example are haloscope experiments that aim to detect axions from the local dark matter halo. They employ highly resonant microwave cavities that enhance the conversion of the axion to the photon in external magnetic fields. These experiments are also sensitive to the dark photon, which mixes with the Standard Model (SM) photon with the practical advantage of not requiring a magnet. However, their sensitivity decreases significantly at masses above 40 μeV .

Therefore, the MAgnetized Disk and Mirror Axion eXperiment (MADMAX) was designed to search for axions in the mass range of 40 μeV to 400 μeV . It uses the dielectric haloscope concept that, instead of a cavity, employs a booster consisting of a stack of dielectric disks and a mirror to enhance the axion-photon conversion. By moving the disks, the axion mass to which the booster is sensitive can be tuned. MADMAX has developed several prototypes categorized as open boosters, which have open boundaries around the disks and mirror, and closed boosters, which are encased in a metal cylinder. To enable free movement of the disks, MADMAX aims at employing an open design for the full-scale setup. The closed booster design allows for easier modeling.

Determining the sensitivity of a dielectric haloscope presents a major challenge. No axion-like calibration source exists and the large detector size compared to the photon wavelength λ of $\mathcal{O}(10^4 \lambda^3)$ renders full-wave simulations difficult. Therefore, a single-mode booster model is presented that is capable of replacing time-consuming and computationally expensive full-wave simulations of closed booster systems. Additionally, connecting a receiver system to the dielectric haloscope modifies its electromagnetic response by reflecting part of the emitted signal power. The change is quantified using a model of the noise emitted by the receiver, which is shown to reproduce power spectra of the dielectric haloscope. The presented methods enabled the first dark matter searches by the MADMAX collaboration.

An axion search with a closed booster prototype, encased in an aluminum cylinder, placed limits on the axion in the mass range 76.56 μeV to 76.82 μeV and 79.31 μeV to 79.53 μeV down to an axion-photon coupling of $g_{a\gamma} \simeq 2 \times 10^{-11} \text{ GeV}^{-1}$. The receiver system caused a decrease of the expected axion signal power by up

to 20 %. The full procedure to determine the experiment’s sensitivity, using the booster and receiver noise model, is presented. The limit is world-leading within its mass range, demonstrating the viability and competitive sensitivity of dielectric haloscopes.

A dark photon search using an open booster prototype, with free-standing disks with open boundaries, placed limits on the dark photon mass range of 78.62 μeV to 83.95 μeV . The peak sensitivity to the kinetic mixing of the dark photon with the SM photon of $\chi \simeq 2 \times 10^{-13}$ surpasses previous limits by up to three orders of magnitude. Its sensitivity in absence of the receiver system was determined as part of the PhD thesis by J. Egge. Using the receiver noise model, the receiver system was found to increase the expected dark photon signal power by up to 25 %, demonstrating that the noise model is applicable to a broad range of configurations.

These dark matter searches provide a solid foundation for future MADMAX experiments. With the necessary tools in place, efforts can now focus on enhancing the experiment’s sensitivity and extending its reach to a broader mass range.

Zusammenfassung

Das Axion und das Dunkle Photon sind zwei viel diskutierte Kandidaten für Dunkle Materie mit erwarteten Massen unterhalb von 1 eV. Das Axion wurde ursprünglich als Lösung für das starke CP-Problem vorhergesagt, d. h. die unerwartete Abwesenheit von Effekten, die die Ladungs-Parität (CP) in der Quantenchromodynamik (QCD) verletzen. Seine schwache Kopplung zum Standardmodell und einfacher Entstehungsmechanismus weckten großes Interesse am Axion als Kandidat für dunkle Materie, und verschiedene Experimente begannen mit der Suche nach ihm. Ein Beispiel sind Haloskop-Experimente, die darauf abzielen, Axionen aus dem lokalen Dunkle-Materie-Halo zu detektieren. Sie setzen dafür hochresonante Mikrowellenresonatoren ein, welche die Umwandlung des Axions in ein Photon in einem externen Magnetfeld verstärken. Diese Experimente sind auch empfindlich gegenüber dem dunklen Photon, das sich mit dem Photon des Standardmodells (SM) mischt, was den praktischen Vorteil hat, dass kein Magnet erforderlich ist. Ihre Empfindlichkeit nimmt jedoch bei Massen über $40\text{ }\mu\text{eV}$ deutlich ab.

Daher wurde das MAgnetized Disk and Mirror Axion eXperiment (MADMAX) entwickelt, welches nach Axionen im Massenbereich von $40\text{ }\mu\text{eV}$ bis $400\text{ }\mu\text{eV}$ sucht. Es nutzt das Konzept des dielektrischen Haloskops, das anstelle eines Resonators einen Booster verwendet, der aus einem Stapel dielektrischer Scheiben und einem Spiegel besteht, um die Umwandlung von Axionen in Photonen zu verstärken. Durch die Bewegung der Scheiben kann die Axionmasse, auf die der Booster anspricht, eingestellt werden. MADMAX hat mehrere Prototypen entwickelt, die in offene Booster, ohne Begrenzungen um die Scheiben und den Spiegel, und geschlossene Booster, die von einem Metallzylinder umschlossen sind, unterteilt werden. Um eine freie Bewegung der Scheiben zu ermöglichen, strebt MADMAX den Einsatz eines offenen Designs für die Anlage in Originalgröße an. Das Design des geschlossenen Boosters ermöglicht eine einfachere Modellierung.

Die Bestimmung der Empfindlichkeit eines dielektrischen Haloskops stellt eine große Herausforderung dar. Es existiert keine axionähnliche Kalibrierungsquelle, und die im Vergleich zur Photonenwellenlänge λ große Detektorgröße von $\mathcal{O}(10^4\lambda^3)$ erschwert Vollwellensimulationen. Daher wird ein single-mode Boostermodell vorgestellt, das in der Lage ist, zeitaufwändige und rechenintensive Vollwellensimulationen geschlossener Boostersysteme zu ersetzen. Zudem verändert der Anschluss eines Empfängersystems an das dielektrische Haloskop dessen elektromagnetischen Eigenschaften, indem ein Teil der emittierten Signalleistung reflektiert wird. Die Veränderung wird anhand eines Modells des vom Empfänger emittierten

Rauschens quantifiziert, welches fähig ist, die Leistungsspektren des dielektrischen Haloskops zu reproduzieren. Die entwickelten Methoden ermöglichten die ersten Suchen nach Dunkler Materie der MADMAX-Kollaboration.

Eine Axion-Suche mit einem geschlossenen Booster-Prototyp, der von einem Aluminiumzylinder umschlossen ist, setzte Limits für das Axion im Massenbereich 76,56 μeV bis 76,82 μeV und 79,31 μeV bis 79,53 μeV bis hinunter zu einer Axion-Photon-Kopplung von $g_{a\gamma} \simeq 2 \cdot 10^{-11} \text{ GeV}^{-1}$. Das Empfängersystem führte zu einer Verringerung der erwarteten Axion-Signalleistung um bis zu 20 %. Das vollständige Verfahren zur Bestimmung der Empfindlichkeit des Aufbaus unter Verwendung des Booster- und Empfänger-Rauschmodells wird vorgestellt. Das Limit ist weltweit führend in diesem Massenbereich und belegt die Durchführbarkeit und wettbewerbsfähige Empfindlichkeit dielektrischer Haloskope.

Eine Suche nach dunklen Photonen unter Verwendung eines offenen Booster-Prototyps mit freistehenden Scheiben resultierte in Limits für einen Massenbereich dunkler Photonen von 78,62 μeV bis 83,95 μeV . Die Spitzenempfindlichkeit für die kinetische Vermischung des dunklen Photons mit dem SM-Photon von $\chi \simeq 2 \cdot 10^{-13}$ übertrifft frühere Limits um bis zu drei Größenordnungen. Die Empfindlichkeit ohne Empfängersystem wurde im Rahmen der Doktorarbeit von J. Egge bestimmt. Unter Verwendung des Empfängerrauschmodells wurde festgestellt, dass das Empfängersystem die erwartete Signalstärke des dunklen Photons um bis zu 25 % erhöht, was zeigt, dass das Rauschmodell auf eine breite Palette von Konfigurationen anwendbar ist.

Diese Suchen nach dunkler Materie bilden eine solide Grundlage für zukünftige MADMAX-Experimente. Da die erforderlichen Werkzeuge nun vorhanden sind, können sich die Bemühungen darauf konzentrieren, die Empfindlichkeit des Experiments zu verbessern und seinen Erfassungsbereich auf einen breiteren Massenbereich auszuweiten.

List of publications

This cumulative thesis consists of three core publications with significant contributions by the thesis author as listed in the list of own contributions:

- A. Ivanov, **D. Leppla-Weber (corresponding author)** et al. (*MADMAX collaboration*), "Sensitivity of a closed dielectric haloscope to axion dark matter" *Accepted by: Journal of Cosmology and Astroparticle Physics* arXiv: 2603.05006 [astro-ph.IM]. url: <https://arxiv.org/abs/2603.05006> [1].
- B. Ary dos Santos Garcia, D. Bergermann, A. Caldwell, V. Dabhi, C. Diaconu, J. Diehl, G. Dvali, J. Egge, E. Garutti, S. Heyminck, F. Hubaut, A. Ivanov, J. Jochum, S. Knirck, M. Kramer, D. Kreikemeyer-Lorenzo, C. Krieger, C. Lee, **D. Leppla-Weber**, X. Li, A. Lindner, B. Majorovits, J.P.A. Maldonado, A. Martini, A. Miyazaki, E. Öz, P. Pralavorio, G. Raffelt, J. Redondo, A. Ringwald, J. Schaffran, A. Schmidt, F. Steffen, C. Strandhagen, I. Usherov, H. Wang and G. Wieching (*MADMAX collaboration*), "First Search for Axion Dark Matter with a MADMAX Prototype" In: *Phys. Rev. Lett.* **135** (4 July 2025), p. 041001. doi: 10.1103/c749-419q. arXiv: 2409.11777. url: <https://link.aps.org/doi/10.1103/c749-419q> [2].
- J. Egge, **D. Leppla-Weber**, S. Knirck et al. (*MADMAX collaboration*) "First Search for Dark Photon Dark Matter with a MADMAX Prototype" In: *Phys. Rev. Lett.* **134** (15 Apr. 2025), p. 151004. doi: 10.1103/PhysRevLett.134.151004. url: <https://link.aps.org/doi/10.1103/PhysRevLett.134.151004> [3].

Additional co-authored publications with minor contributions by the thesis author:

- J. Egge, M. Ekmedžić, A. Gardikiotis, E. Garutti, S. Heyminck, C. Kasemann, S. Knirck, M. Kramer, C. Krieger, **D. Leppla-Weber**, S. Martens, E. Öz, N. Salama, A. Schmidt, H. Wang and G. Wieching, "Experimental determination of axion signal power of dish antennas and dielectric haloscopes using the reciprocity approach". In: *Journal of Cosmology and Astroparticle Physics* **2024.04** (Apr. 2024), p. 005. issn: 1475-7516. doi: 10.1088/1475-7516/2024/04/005. url: <https://iopscience.iop.org/article/10.1088/1475-7516/2024/04/005> [4].
- M. A. Akgümüş, N. Salama, J. Egge, E. Garutti, M. Maroudas, L. H. Nguyen, **D. Leppla-Weber**, "A new Limit for Axion Dark Matter with SPACE", Accepted by: *Journal of Cosmology and Astroparticle Physics*, arXiv: 2506.18411 [hep-ex]. url: <https://arxiv.org/abs/2506.18411> [5].
- S. Spannagel, S. Lachnit, H. Perrey, J. Braach, L. Kristian Bryngemark, E. Garutti, A. Herkert, F. King, C. Krieger, **D. Leppla-Weber**, L. Ros, S. R. Daza, M. Safdari, L. G. Sarmiento, A. Vauth, H. Wennlöf, "Constellation: The autonomous control and data acquisition system for dynamic experimental setups" In: *Nucl. Instr. Meth. A* **1085** (2026), p. 171279. issn: 0168-9002. doi: 10.1016/j.nima.2026.171279. arXiv: 2601.06494 [physics.ins-det] [6].
- D. Maglione, S. Knirck, **D. Leppla-Weber**, "Axion Dark Matter EXperiment 2A Cavity Characterization" In: *FNAL report* [7]

Contents

List of publications	vii
1 Introduction	1
2 Wave-like dark matter	5
2.1 Axion dark matter	6
2.1.1 QCD axion	7
2.1.2 Axion cosmology	8
2.1.3 Axion-photon coupling	10
2.1.4 Axion-like particles	11
2.2 Dark photon dark matter	12
3 Searching for axion dark matter	13
3.1 Haloscope searches	15
3.1.1 Dish antenna	17
3.1.2 Cavity haloscopes	17
3.2 Dielectric haloscopes	18
3.2.1 The Magnetized Disk and Mirror Axion Experiment	22
3.3 Relation to dark photon searches	23
4 Receiver system for a dielectric haloscope	24
4.1 Dark matter signal	25
4.2 Heterodyne detection	26
4.3 Signal-to-noise ratio	27
4.4 Suppression of unwanted signals	31
4.4.1 Signals in the RF band	31
4.4.2 Signals in the IF band	32
4.4.3 Signals due to a non-ideal receiver system	34
4.5 Amplifier mismatch	35
5 Sensitivity of a closed dielectric haloscope to axion dark matter	36
5.1 Introduction	36
5.2 Booster and receiver noise model	39
5.2.1 Booster model	40
5.2.2 Receiver noise model	44
5.2.3 Combined noise model	47

5.3	Model verification	48
5.3.1	Ideal case	49
5.3.2	Treatment of 3D effects	49
5.3.3	Control of higher-order modes	55
5.3.4	Identification of the booster mode	56
5.4	Sensitivity of a closed booster axion search	58
5.4.1	Receiver chain characterization	59
5.4.2	Boost factor determination	59
5.4.3	Uncertainty propagation	63
5.5	Conclusion	64
5.6	Appendix	66
5.6.1	Receiver power calibration	66
5.6.2	Model parameters	68
5.7	Additional material	70
5.7.1	Analytical expression of the noise model	70
6	Search for axion dark matter using a closed booster	73
6.1	Introduction	73
6.2	Experimental Setup	75
6.3	Boost factor determination	76
6.4	Analysis	77
6.5	Conclusion	79
6.6	Additional material	82
6.6.1	Monitoring system	82
6.6.2	Effect of the receiver system on the boost factor	82
7	Search for dark photon dark matter using an open booster	86
7.1	Introduction	86
7.2	Experimental Setup	88
7.3	Data Taking Run	92
7.4	Analysis	93
7.5	Conclusion	95
7.6	Additional material	96
7.6.1	Bin masking and frequency hopping	96
7.6.2	Power calibration	97
7.6.3	Application of the noise model	97
8	Conclusion	101
	List of own contributions	104
	Acknowledgments	106

1 Introduction

The nature of dark matter (DM) is one of the greatest mysteries in modern physics. It is assumed to dominate the matter content of the universe, but so far, only its gravitational effects have been observed and the exact nature remains unknown. Evidence for its existence ranges from the scale of the entire universe, for example, effects on the cosmic microwave background (CMB) from structure formation [8], to the scale of galaxies, for example, their unexpected velocity distributions [9]. As discussed at the beginning of chapter 2, the observations lead to the conclusion that DM is made of non-SM particles with negligible couplings to the SM and has to be non-relativistic. From the observed temperature anisotropies of the CMB, the DM energy density in the universe is estimated to be 26 %, while ordinary baryonic matter only constitutes about 5 % [8]. The remaining 69 % are attributed to dark energy. Measuring the velocity distribution of the astrophysical objects within different galaxies shows a significant discrepancy between the velocities required for the objects to remain gravitationally bound and the observed velocities. From the measured difference in the Milky Way, the local DM density ρ is estimated to be $300 \text{ MeV/cm}^3 \lesssim \rho \lesssim 600 \text{ MeV/cm}^3$ [10]. Despite the large abundance, DM is continuing to evade detection.

There are various theories about what might constitute DM, spanning over 50 orders of magnitude in mass. A brief overview is provided at the beginning of chapter 2. This cumulative thesis focuses on the experimental search for axion and dark photon DM with masses around $100 \mu\text{eV}$. The axion in particular has gained significant prominence as a DM candidate in recent decades and is therefore introduced more thoroughly in section 2.1. It is a consequence of a proposed solution to the strong CP problem [11], that is, the unexpected absence of charge-parity (CP) violation within quantum-chromodynamics (QCD), and was only later realized to be an excellent DM candidate [12]. Its mass was originally predicted to be of $\mathcal{O}(100 \text{ keV})$ [12], which was quickly ruled out experimentally [13]. Further astrophysical constraints exclude the axion at masses $\gtrsim 0.1 \text{ eV}$ [14]. Assuming it is responsible for the entirety of the observed DM, some simulations predict axion masses around $100 \mu\text{eV}$ [15, 16], as considered in this thesis. In addition to the axion, the dark photon provides another DM candidate that can be probed using similar experimental setups. It arises as a massive vector boson in minimal extensions to the Standard Model (SM) and could in principle provide a portal to a hidden DM sector [17].

Existing direct searches [18, 19, 20] based on nucleon scattering are not sensitive to axions at masses below 0.1 eV, since they transfer too little momentum to the detector. Furthermore, the expected couplings to the SM are too small to lead to significant production in accelerator-based searches. Thus, dedicated experiments have been designed that can be categorized as light-shining-through-wall (LSW) experiments [21], helioscopes [22] and haloscopes [22]. These experiments rely on the conversion of axions to SM photons in the presence of magnetic fields. They differ significantly in terms of the axion source. LSW experiments are purely laboratory-based and produce their own axions using a laser in a magnetic field. Helioscopes aim to detect axions produced in the Sun, while haloscopes search for axions from the galactic DM halo. These experiments are also generally sensitive to dark photons, which have the practical advantage of not requiring a magnetic field. This thesis focuses on haloscope experiments, which have the highest sensitivities but are also dependent on the rather uncertain determination of the local DM density.

Haloscopes, introduced in chapter 3, typically use highly resonant microwave cavities that enhance the axion-photon conversion. The resulting signal power can be increased by stronger resonances and larger cavity volumes. However, the volume of the cavity scales inversely with the photon frequency, which, in the non-relativistic limit, is set by the axion mass as $\nu = m_a c^2 / h$. Therefore, it is challenging to reach sufficient sensitivity to masses greater than 40 μeV with such experiments.

To overcome this limitation, the Magnetized Disk and Mirror eXperiment (MADMAX) working group, which led to the founding of the MADMAX collaboration, proposed the dielectric haloscope [23]. Instead of relying on a resonant cavity, it uses a stack of dielectric disks and a mirror, called *booster*, to enhance the axion-photon conversion in a magnetic field. This results in coherent photon emissions at the interfaces, as explained in more detail in section 3.2. The dielectric haloscope’s sensitivity can be improved by increasing the radius and number of disks. Both numbers can be scaled independently of the probed DM mass, which allows searches for DM at higher masses than cavity experiments.

MADMAX plans to build a dielectric haloscope consisting of up to 80 disks with a surface area of $\sim 1 \text{ m}^2$ together with a 9 T dipole magnet to probe axion masses from 40 μeV to 400 μeV [24]. On its way to achieve this goal, MADMAX has been working on an extensive prototyping program. To study the electromagnetic properties of a dielectric haloscope, a first proof-of-principle setup has been built [25], consisting of up to five free-standing sapphire disks and a mirror. Such a setup without boundaries around the booster is called an *open booster*. The signal power emitted by the booster is typically focused onto a horn antenna using a focusing mirror. It has been used to study basic tuning capabilities and

mechanical requirements. As a result, a system to move the disks with $\mathcal{O}(\mu\text{m})$ precision at cryogenic temperatures and in a strong magnetic field was developed and tested [26]. To ease the modeling of the booster, MADMAX has developed so-called *closed booster* systems, which are encased in a metal cylinder. Unlike the open design, the closed boundary conditions result in a finite set of propagating modes within the system, which enables the use of simplified models.

The central goal of this cumulative thesis is to develop a method to determine the sensitivity of a dielectric haloscope to subsequently search for DM. This involves two key challenges: the simulation of the large booster volume and the quantification of the previously unaccounted-for effect of the receiver system on the DM signal power.

Simulations are required, because the photon emissions at the disks and mirror due to the axion-photon conversion are difficult to reproduce experimentally. A direct calibration with an axion-like source is therefore not possible. However, full-wave simulations are also challenging because of the large detector volume. Simulating the full system requires a domain size exceeding $\mathcal{O}(10^4\lambda^3)$ compared to the photon wavelength λ . This is orders of magnitude larger than a typical cavity haloscope with a size of $\mathcal{O}(10\lambda^3)$.

Furthermore, the booster's response to the axion is directly modified by the presence of a receiver system, as described in chapter 4. The expected signal power emitted by the booster is typically below $\mathcal{O}(10^{-21}\text{ W})$, necessitating highly sensitive amplifier devices. However, these amplifiers modify the electromagnetic response of the dielectric haloscope by reflecting part of the potential signal. This problem is typically addressed by adding a device that is only transmissive in one direction, such as a circulator, in front of the receiver system. However, such devices cannot be operated in magnetic fields and decrease the experimental sensitivity, especially at higher frequencies.

The first publication included in this cumulative thesis in chapter 5 addresses both of these challenges. It presents a method to determine the sensitivity of a closed booster to axion DM, including the effect of the receiver system. The procedure is based on a single-mode model, which is validated using full-wave finite-element method (FEM) simulations and is shown to accurately describe the sensitivity of the booster in the absence of the receiver system. This solves the issue of computationally expensive simulations for closed boosters. The publication further includes a model of the noise emitted by the receiver system that is capable of reproducing power spectra of the booster. Using this noise model, the effect of the receiver system on the sensitivity of the experiment can be quantified. Sensitivity-decreasing devices that mitigate this effect are therefore no longer necessary.

Consequently, an axion DM search was performed using a closed MADMAX

prototype called CB200 in the 1.6 T Morpurgo magnet at CERN. Its sensitivity is determined using the method introduced above. The results are presented in the second publication included in this thesis in chapter 6.

MADMAX aims at developing an open booster to be able to freely tune the position of the disks. Such a system is no longer sufficiently described by a single-mode model. To study the properties of this design, MADMAX has built a three-disk open booster called OB300. The sensitivity determination in the absence of the receiver system was solved in the PhD thesis by J. Egge [27] using a measurement-based approach called the bead-pull method [4]. Since no magnet was available, it was used to search for dark photon DM. The main contribution in the context of this thesis is the implementation and characterization of the receiver system, including its effect on the experiment's sensitivity. The publication containing the results of the dark photon search is included in chapter 7, with additional unpublished material produced as part of this thesis shown in section 7.6.

In summary, this thesis makes three significant contributions to the development of dielectric haloscopes: it develops a computationally efficient single-mode model that replaces full-wave simulations for closed booster systems, it introduces a receiver noise model that quantifies the impact of the receiver system on the experiment's sensitivity and it applies these methods to the first two DM searches with MADMAX prototypes.

2 Wave-like dark matter

Despite the overwhelming evidence, the exact nature of DM remains unknown. The first hints for its existence came from observations of the rotation curves of spiral galaxies and the dynamics of galaxy clusters.

The mass contained in galaxy clusters can be estimated from the observed velocity dispersion. This was first done by Zwicky [28], who applied the virial theorem to the Coma cluster to relate the average kinetic to the average potential energy. This method yielded a significantly higher mass than what was visible, again indicating some form of invisible mass. The fact that DM seems to be gravitationally bound in such large-scale structures, already suggests that it cannot be relativistic.

Spiral galaxies have their visible mass mostly concentrated within $\mathcal{O}(10\text{ kpc})$ of the galactic center. At larger radii r , the circular velocity of astrophysical objects is expected to decrease as $r^{-1/2}$ to remain gravitationally bound. However, observations show that the circular velocity remains roughly constant at large radii, suggesting the presence of additional, invisible mass in the form of a DM halo [9].

The most quantitatively robust constraints on the properties of DM arise from its role in the formation of large-scale structures in the early universe. In the absence of DM, structures such as the galaxies observed today could not have formed. Radiation pressure would have prevented baryonic matter from collapsing under its own gravity, making an additional, non-interacting source of gravitational attraction necessary. Consequently, DM has to be *cold*, i.e. non-relativistic, and cannot be a SM particle, since it otherwise would have been exposed to the same radiation pressure. The process of structure formation leaves imprints on the CMB, that can be reproduced with high accuracy by the Λ CDM model [8]. Here, CDM stands for cold dark matter and Λ represents the cosmological constant, which is associated with dark energy. The results indicate that 26 % of the energy density in the universe consists of non-SM cold DM.

Because only gravitational effects have been observed, modifications to gravity have been considered instead of inventing a new type of matter [29]. However, the existence of DM is the only theory so far that is capable of consistently explaining all of the observations [30].

A wide variety of theories on the nature of DM has been proposed, ranging from primordial black holes (PBHs) heavier than the Sun to ultralight scalar fields with masses below 1 eV. A comprehensive review can be found at [30]. PBHs,

in particular, represent an attempt to account for DM using baryonic matter. They are hypothesized to have formed before the epoch of large-scale structure formation, allowing them to evade early universe constraints. However, they have been mostly excluded by observations of gravitational waves and gravitational lensing effects [31].

In the previous decades, significant effort was put into searching for Weakly Interacting Massive Particles (WIMPs) [32]. Such particles are commonly predicted by supersymmetric extensions of the SM and are expected to be of masses of $\mathcal{O}(\text{GeV})$. Many direct searches [18, 19, 20] have excluded large regions of the WIMP parameter space, reaching the so-called neutrino fog where the SM neutrino becomes the dominant background of the experiment.

More recently, wave-like DM has emerged as a topic of interest. The name is motivated by the fact that such particles with masses $< 1 \text{ eV}$ are best described by classical fields, due to their high number density relative to their de Broglie wavelength. From the same reasoning, they are required to be bosonic to not violate the Pauli exclusion principle. Ultralight bosonic fields are a common prediction by various theories [33]. Most current searches are, however, not sensitive to this parameter space. Their mass is below the reach of current direct searches and the expected coupling is too small to enable significant production in accelerator-based experiments. The strong theoretical motivation, together with the lack of sensitivity of current experiments, motivates the usage of new detector concepts.

Before going into the details of the detection in chapter 3, this chapter introduces the two wave-like dark matter candidates considered in this work. The axion is the main motivation of the presented experiment and is therefore introduced in more detail in section 2.1. Section 2.2 is then dedicated to briefly describing the dark photon.

2.1 Axion dark matter

The axion is a popular wave-like DM candidate due to its theoretical origin from a solution to the strong CP problem, presented in section 2.1.1. It has all the properties required of DM, namely a small interaction with the SM and sufficient, non-relativistic, production during the early universe to constitute all of the observed DM. In section 2.1.2 different cosmological scenarios are described including their effect on the prediction of the axion mass. Section 2.1.3 then introduces the coupling of axions to the SM photon. In section 2.1.4, axion like particles (ALPs) are discussed briefly.

2.1.1 QCD axion

The axion is the consequence of a solution to the strong CP problem proposed by Peccei and Quinn [11]. The strong CP problem describes the observed absence of CP violating effects in the QCD sector, although the QCD Lagrangian contains a term that is in principle CP violating [34]:

$$\mathcal{L} = \frac{g_s}{32\pi} \theta G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}, \quad (2.1)$$

with the strong coupling constant g_s , the gluonic field strength tensor G and its dual $\tilde{G}$ and a free angular parameter θ that can in principle range from $-\pi$ to π . The angle θ , containing only QCD interactions, is however not a physical observable. One needs to also consider effects from weak interactions, that are quantified by the quark mixing matrix M . This leads to the observable effective angle $\bar{\theta} = \theta + \arg \det M$ [34].

A non-zero $\bar{\theta}$ would lead to a non-zero electric dipole moment (EDM) of the neutron given approximately by $d_n = 8.2 \times 10^{-16} \bar{\theta} \text{ e cm}$ [35]. Current measurements constrain the neutron EDM to be $d_n < 1.8 \times 10^{-26} \text{ e cm}$ [36] and therefore require $\bar{\theta} < 10^{-10}$. Because $\bar{\theta}$ contains contributions from the strong and the electroweak sector, that are both expected to be of $\mathcal{O}(1)$, $\bar{\theta} = 0$ requires what is considered unnatural fine tuning.

Peccei and Quinn proposed to promote $\bar{\theta}$ to a dynamical field $\bar{\theta} \rightarrow a(\mathbf{x}, t)/f_a$ by introducing an additional global $U_{\text{PQ}}(1)$ symmetry that is spontaneously broken at the energy scale f_a [11]. The QCD potential is then naturally minimized at $\bar{\theta} = 0$, explaining the observed absence of CP violation. Excitations of the new field $a(\mathbf{x}, t)$ give rise to a new pseudo-scalar boson. This was only later realized by Wilczek and Weinberg [12], who also gave it the name *axion*.

In its most general form, the axion only interacts with gluons at the tree level. However, it has identical quantum numbers as the neutral pion π^0 , i.e. zero electric charge, spin-0 and odd-parity. By first mixing with the π^0 , the axion effectively couples to two photons. This provides the most commonly used detection channel and is explored further in section 2.1.3. The axion also generally couples to nucleons and, depending on the model, can couple to electrons at the tree-level [37]. The nucleon coupling is rarely used for axion searches, with CASPER [38] being a notable exception. If the axion couples to electrons, this could potentially lead to an increase of the axion flux from the Sun by more than a factor of 50, which would drastically increase the sensitivity of helioscope experiments [39].

The axion's mass is linearly related to the energy scale f_a and from its mixing with pions was calculated to be [40]

$$m_a \simeq 5.7 \text{ peV} \left(\frac{10^{12} \text{ GeV}}{f_a} \right). \quad (2.2)$$

It was originally predicted to be of the order of $m_a \simeq 100 \text{ keV}$ [12], by assuming $f_a \simeq 250 \text{ GeV}$ equivalent to the electroweak scale. This was however quickly excluded experimentally [13].

In principle, f_a can be arbitrarily large such that the axion mass itself becomes small enough to evade immediate detection. This realization led to the development of so-called *invisible axion* theories. The most commonly referenced theories are Kim-Shifman-Vainshtein-Zakharov (KSVZ) type theories [41, 42], that introduce additional heavy quarks, and Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) type theories [43], that introduce an additional Higgs doublet. To distinguish the axion that was introduced to solve the strong CP problem from phenomenologically similar particles predicted by other theories, it is commonly referred to as the QCD axion.

2.1.2 Axion cosmology

The axion is produced during the early universe by the so-called misalignment mechanism. The $U_{\text{PQ}}(1)$ symmetry is spontaneously broken at the energy scale f_a , which sets an initial misalignment angle θ_i and results in a complex scalar field $\Phi = |\Phi|e^{i\theta}$. The phase θ is interpreted as the axion, which in principle has a shift symmetry $\theta \rightarrow \theta + \alpha$. This symmetry is explicitly broken at the Lagrangian level at a temperature of $T_{\text{QCD}} \simeq \mathcal{O}(100 \text{ MeV})$ by the coupling to gluons, therefore the axion acquires a mass term [44]. The initial misalignment from the potential minimum leads to field oscillations, interpreted as a set of axion particles.

One crucial requirement of DM is that it scales as matter during the expansion of the universe, that is its density $\rho_{\text{DM}} \propto a^{-3}$ with the scaling factor a . By solving the equation of motion of the axion scalar field, it can be shown that this is naturally true if $m_a(T_{\text{eq}}) > 3H(T_{\text{eq}})$, with the axion mass m_a and Hubble radius H at the time of matter-radiation equality T_{eq} [44]. This sets a lower bound on the axion mass of approximately 10^{-28} eV [44]. The resulting axion energy density today Ω_a , produced by the misalignment mechanism, can be written in terms of the initial misalignment angle θ_i and the axion mass m_a as [44]

$$\Omega_a h^2 \simeq 0.12 \theta_i^2 \left(\frac{4.7 \text{ } \mu\text{eV}}{m_a} \right)^{\frac{7}{6}}. \quad (2.3)$$

The normalization to 0.12 is chosen to reflect the DM density in the universe as determined from the CMB. Assuming the axion provides all of DM, the axion mass can therefore be written as

$$m_a = 4.7 \text{ } \mu\text{eV} \theta_i^{12/7}. \quad (2.4)$$

During the early universe, the space is expected to expand exponentially referred to as inflation. Whether the $U(1)_{\text{pq}}$ symmetry is broken before or after inflation

takes place has a big impact on the predictability of θ_i and therefore also m_a . In the pre-inflationary case, the symmetry breaking results in a constant, but randomly chosen θ_i for the observable universe. The axion mass is therefore fundamentally unpredictable and only constrained by the full range of $|\theta_i| \leq \pi$, although one might make a naturalness argument for $\theta_i \simeq 1$. Assuming the axion provides all of the observed DM, this would then result in an axion mass around $5 \mu\text{eV}$ as per equation (2.4).

In the post-inflationary case, the universe consists of causally disconnected patches at the time of the symmetry breaking. This leads to a different θ_i angle in each patch. The axion mass, assuming the axion is responsible for all of the observed DM, can then in principle be predicted by an average over all θ_i angles to be approximately $26 \mu\text{eV}$ [45]. However, this case also leads to additional axion production mechanisms from topological defects such as cosmic strings and domain walls. The cosmic strings are a result of the phase θ , corresponding to the axion, creating loops from $-\pi$ to π due to the random initial misalignment θ_i . The complex scalar field Φ is then artificially forced to 0 at the center, which is a potential maximum. These singularities form along a line, called a cosmic string. Cosmic strings form before the QCD cross-over, therefore the axion does not have its preferred vacuum expectation value of $\theta = 0$ yet. After this transition, the strings can unravel with θ relaxing towards 0. However, this can result in two-dimensional sheets between strings, where the axion is stuck at a saddle point of the potential. Those are called domain walls. The energy stored in those singularities leads to additional axions being radiated.

Additional axion production mechanisms result in less production from the misalignment mechanism to provide the required abundance of DM. As apparent from equation (2.3), this necessitates a higher axion mass. Therefore $m_a \simeq 26 \mu\text{eV}$ presents a rather strict lower bound in this scenario. Many attempts at simulating the axion production in the post-inflationary scenario exist, that can be roughly summarized as a prediction of $26 \mu\text{eV} \lesssim m_a \lesssim 500 \mu\text{eV}$ [44] for the axion to provide the entirety of the observed DM.

Having additional axion production after the time of inflation causes local overdensities of axions that collapse to so-called axion miniclusters [46]. This has some consequences for axion detection: about 75 % of axion DM is expected to be bound in clumps of the mass of an asteroid and the size similar to the Earth's orbit [47]. It is therefore likely that a detector will be looking at the space between such miniclusters that might have an axion density as little as 10 % of the estimated galactic DM density [47]. This value is however to be seen as a conservative lower bound.

2.1.3 Axion-photon coupling

The expected DM density of $\rho \simeq 300 \text{ MeV/cm}^3$ together with the low masses $< 1 \text{ eV}$ lead to a very high axion number density $n_a > 3 \times 10^8 \text{ cm}^{-3}$. In combination with the large de Broglie wavelength due to the non-relativistic nature of CDM, this allows for a treatment of the axion as a classical field $a(\mathbf{x}, t)$. The effective Lagrangian, describing the interaction of axions with the electric field $\mathbf{E}$ and magnetic field $\mathbf{B}$ for a given axion-photon coupling $g_{a\gamma}$, then reads:

$$\mathcal{L} = g_{a\gamma} a \mathbf{E} \cdot \mathbf{B}. \quad (2.5)$$

This leads to a modification of Maxwell's equations, specifically Ampere's law [48]

$$\nabla \times \mathbf{H} - \dot{\mathbf{D}} = \mathbf{J} + g_{a\gamma} (\mathbf{B} \dot{a} - \mathbf{E} \times \nabla a), \quad (2.6)$$

with the magnetic field $\mathbf{H} = \mathbf{B} - \mathbf{M}$, displacement field $\mathbf{D} = \mathbf{E} + \mathbf{P}$ and current density $\mathbf{J}$. $\mathbf{M}$ and $\mathbf{P}$ absorb the effects of magnetization and polarization, respectively. The non-linear extra term, containing products of a with $\mathbf{B}$ and $\mathbf{E}$, presents the most common detection channel of the axion. In the presence of a magnetic field $\mathbf{B}$ or electric field $\mathbf{E}$, the axion field acts as an additional source term. The axion-sourced fields can again interact with the axion, making the term non-linear. Experimentalists typically aim to maximize the source term by the use of strong external fields. In this case, higher-order terms can be neglected and the fields within equation (2.6) replaced by their external components $\mathbf{E}_e$ and $\mathbf{B}_e$. The full linearized Maxwell's equations then read [48]:

$$\nabla \cdot \mathbf{B} = 0 \quad (2.7a)$$

$$\nabla \times \mathbf{E} = -\dot{\mathbf{B}} \quad (2.7b)$$

$$\nabla \cdot \mathbf{D} - \rho_f = -g_{a\gamma} \mathbf{B}_e \cdot \nabla a \quad (2.7c)$$

$$\nabla \times \mathbf{H} - \dot{\mathbf{D}} = \mathbf{J} + g_{a\gamma} (\mathbf{B}_e \dot{a} - \mathbf{E}_e \times \nabla a). \quad (2.7d)$$

The additional source term is commonly defined as the effective axion current:

$$\mathbf{J}_a \equiv g_{a\gamma} (\mathbf{B}_e \dot{a} - \mathbf{E}_e \times \nabla a). \quad (2.8)$$

Solving the linearized axion modified Maxwell's equations (equation (2.7)) for a given set of boundary conditions with only $\mathbf{J}_a$ as a source leads to the axion-induced electric and magnetic fields $\mathbf{E}_a$ and $\mathbf{B}_a$ that can be detected by experiments. In a vacuum and assuming an axion field of constant amplitude $a(t) = a_0 e^{im_a t}$, corresponding to the non-relativistic limit, the axion induced field is given by

$$\mathbf{E}_a(t) = -g_{a\gamma} \mathbf{B}_e a_0 e^{im_a t}. \quad (2.9)$$

The amplitude of the axion field can be further expressed in terms of the axion density ρ_a as [48]

$$a_0 = \frac{\sqrt{2\rho_a}}{m_a}. \quad (2.10)$$

Given a specific axion model, the expected axion-photon coupling $g_{a\gamma}$ is related to the axion mass m_a as [49]

$$g_{a\gamma} \simeq 2 \times 10^{-14} \text{ GeV}^{-1} \left(\frac{m_a}{100 \mu\text{eV}} \right) C_{a\gamma}, \quad (2.11)$$

where $C_{a\gamma}$ is a model dependent, dimensionless constant. The prediction of the DFSZ ($C_{a\gamma} = 0.75$) and KSVZ ($C_{a\gamma} = -1.92$) [49] models is often taken as a benchmark for the sensitivity of axion searches and is also the main motivation of the experiment presented in this work. However, considering for example the existence of magnetic monopoles [50], the expected axion-photon coupling can be increased by approximately five orders of magnitude.

2.1.4 Axion-like particles

Axion-like particles (ALPs) have the same phenomenology as the QCD axion but do not couple to gluons at the tree-level. They therefore no longer solve the strong CP problem and have no direct connection between the axion mass m_a and the energy scale f_a (equation (2.2)).

ALPs typically either arise as pseudo-Nambu-Goldstone bosons due to a broken global symmetry or as zero modes of antisymmetric tensor fields during compactification in string theory [49, 33]. The exact number of ALPs predicted by a specific compactification varies, with the distribution for known Calabi-Yau manifolds peaking at ~ 30 [49]. Just like the QCD axion, ALPs can be produced abundantly during the early universe from the misalignment mechanism and therefore are also a good DM candidate [33].

Apart from the standard misalignment mechanism, one might also consider the so-called kinetic misalignment mechanism [51], where the axion field has a non-zero initial velocity. Together with additional couplings to the SM, this can potentially lead to the observed baryon – antibaryon asymmetry of our universe by what is called ALPogenesis [52].

Strictly speaking, one has to differentiate between QCD axions, that couple to gluons and solve the strong CP problem and ALPs that do not. From a detection point of view there is however no conceptual difference between the two, hence this work refers to both, the QCD axion and ALPs, when mentioning axions.

2.2 Dark photon dark matter

The dark photon is introduced to the SM by an additional $U(1)_\chi$ gauge symmetry, that is broken during the early universe [17]. This results in a massive spin-1 vector boson. The dark photon is typically motivated as a gauge boson that mediates interactions between the SM and some hypothetical dark sector. If other SM particles are uncharged under the $U(1)_\chi$ symmetry, the dark photon only interacts with them gravitationally and therefore also makes for a good DM candidate. Its quantum numbers are identical to the SM photon, therefore they mix kinetically. This leads to a modification of Ampere's law [53], similar to the one introduced due to the axion:

$$\nabla \times \mathbf{H} - \dot{\mathbf{D}} = \mathbf{J}_f + \chi(\dot{\mathbf{E}}_\chi - \nabla \times \mathbf{B}_\chi) \quad (2.12)$$

with the kinetic mixing angle χ , the dark photon electric field $\mathbf{E}_\chi$ and the dark photon magnetic field $\mathbf{B}_\chi$. Analogous to the effective axion current, one then defines the effective dark photon current:

$$\mathbf{J}_{dp} \equiv \chi(\dot{\mathbf{E}}_\chi - \nabla \times \mathbf{B}_\chi). \quad (2.13)$$

All equations derived from $\mathbf{J}_a$ consequently also apply to $\mathbf{J}_{dp}$. A major difference from the effective axion current is that $\mathbf{J}_{dp}$ does not depend on an external magnetic field. Therefore, no (typically expensive) magnet is required for detection. It is also important to note, that $\mathbf{E}_\chi$ is a vector field with unknown polarization. Detectors are often only sensitive to a single polarization. For a fixed DP polarization, this would therefore require a specific orientation between the experiment and the dark photon field $\mathbf{E}_\chi$. Instead, a random polarization is commonly assumed, such that a statistical average over all possible polarizations can be used to simplify the analysis. A thorough review of the effect of the dark photon polarization on haloscope searches can be found at [54].

3 Searching for axion dark matter

Three types of experiments have been specifically developed to search for the axion: Light-shining-through-the-wall (LSW) experiments, helioscopes, and haloscopes.

LSW experiments [21] aim to detect axions that are generated in the laboratory by a strong laser within a magnetic field and therefore place the fewest assumptions on the production mechanism of axions. They provide a general way of searching for pseudo-scalar and scalar particles at masses $\lesssim 1$ eV instead of specifically searching for axion DM. Due to the inverse relation between mass and energy scale, such experiments can probe very high energies $> 10^{10}$ GeV that are inaccessible to accelerator physics. This makes them an important addition to particle physics searches. Their sensitivity is, however, limited, since they rely on the successive conversion of photons into axions and back into photons. Therefore, they are less likely to find the QCD axion, although some theories [50, 52] also motivate couplings accessible by LSW experiments like ALPS II [55].

Helioscopes [22] instead point a telescope containing a magnet and an X-ray sensor at the sun, with the goal of detecting relativistic axions produced there. The high photon density within the sun together with strong electric fields from charged particles in the hot plasma result in an expected high axion flux towards Earth of $\mathcal{O}(10^9 \text{ cm}^{-2} \text{ sec}^{-1})$ [22] for $g_{a\gamma} \simeq 10^{-12} \text{ GeV}^{-1}$. This allows helioscopes to reach higher sensitivities than LSW experiments. One example is the CAST experiment [56], which has placed limits down to $g_{a\gamma} \simeq 5.7 \times 10^{-11} \text{ GeV}^{-1}$ over a broad axion mass range as visible in figure 3.1. Future searches by (Baby)IAXO [57] plan to increase the limits placed by CAST by about 1.5 orders of magnitude. Consequently, it is expected to reach sensitivity to the QCD axion at masses $\gtrsim 1$ meV.

So far, only haloscope experiments [22] have been able to reach sensitivity to the axion-photon coupling predicted by the KSVZ and DFSZ models at masses $\lesssim 10$ meV. Multiple haloscopes exist [58, 59, 60, 61] that use resonant microwave cavities to detect axions from the DM halo. They can achieve high sensitivities, but typically only cover a narrow mass range. Furthermore, they rely on the rather uncertain determination of the local DM density and assume it is fully provided by the axion.

In figure 3.1, the current exclusions on the axion-photon coupling $g_{a\gamma}$ are summarized for axion masses $10^{-12} \text{ eV} \leq m_a \leq 10 \text{ eV}$. Astrophysical observations are shown in green, exclusions by helioscopes in dark red and exclusions by haloscopes in bright red, assuming a DM density of $\rho = 450 \text{ MeV cm}^{-3}$. The axion-photon

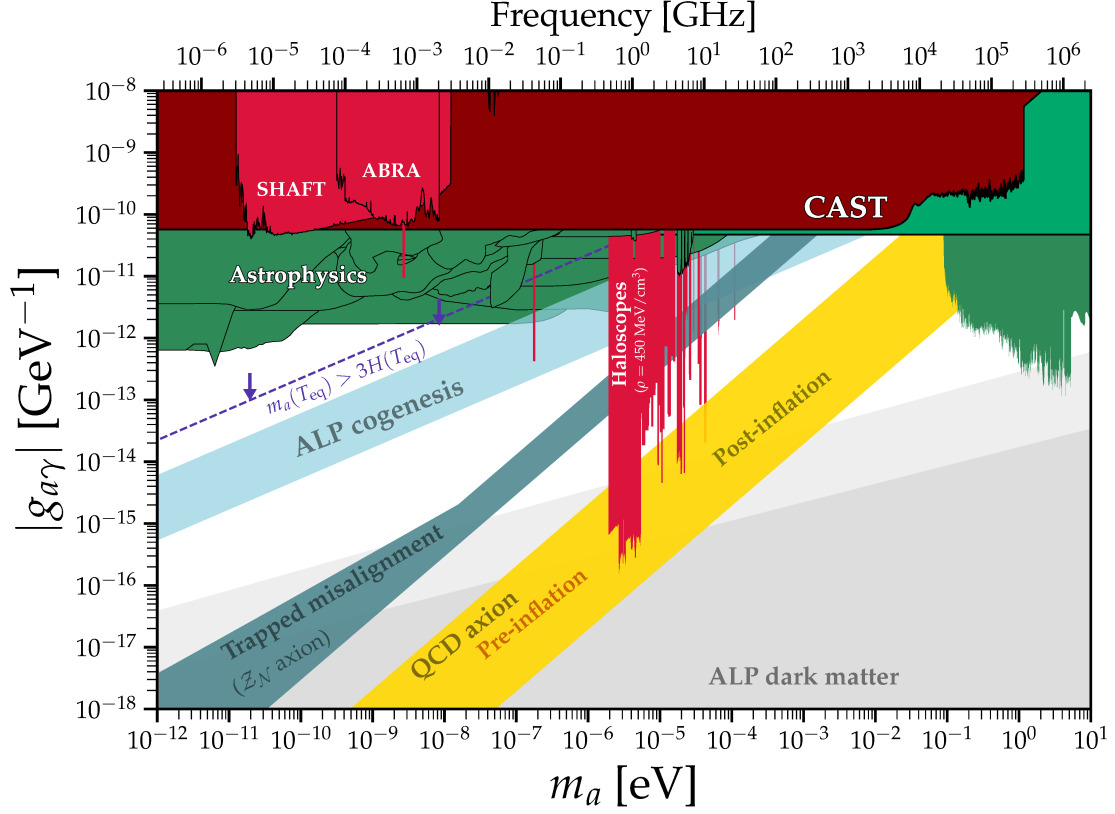


Figure 3.1: Existing exclusion limits on the axion parameter space. Exclusions from astrophysical observations are shown in green, from helioscopes in dark red and from haloscopes in bright red. The axion-photon coupling predicted by the KSVZ and DFSZ models is indicated by the yellow band. Additional predictions are shown for the cogenesis [52] and trapped misalignment [62] cases. The expected parameter space of ALP DM in the case that the entirety of DM is produced by the standard misalignment mechanism is indicated in gray. Axions with couplings above the dashed blue line are excluded as DM since the axion field’s density would not scale as matter during the evolution of the universe. Produced with *AxionLimits* by C. O’Hare [63].

coupling predicted by the DFSZ and KSVZ theories is indicated by the yellow band. A general ALP model where the entirety of DM is produced with the standard misalignment mechanism results in an allowed parameter space indicated by the gray area [33]. The different shades of gray correspond to models with a constant axion mass during the expansion of the universe (dark gray) or a non-constant mass (light gray). By a modified misalignment mechanism, the expected coupling can be increased significantly, as shown for the ALPogenesis [52] and trapped misalignment [62] cases.

A large fraction of the parameter space predicted by simulations of the post-inflationary scenario remains unexplored, providing the main motivation of this thesis. The sensitivity of haloscopes to the axion primarily depends on the signal power that is extracted from the axion field by the experimental setup. Therefore, in section 3.1 the signal power achievable by existing haloscope experiments in this mass range is derived and shown to be insufficient to reach sensitivity to this parameter space. This motivates the dielectric haloscope concept, which is introduced together with its challenges in section 3.2.

3.1 Haloscope searches

Haloscope experiments aim to detect axions from the galactic DM halo by exploiting the axion-photon conversion in magnetic fields. Their setup can generally be split into two parts, the optical system and the receiver system. The optical system is responsible for converting axions to photons that are subsequently detected by the receiver, which will be introduced in chapter 4. This section focuses on determining the signal power emitted by the optical system. Subsequently, the achievable signal power of dish antenna experiments and cavity haloscopes is estimated in section 3.1.1 and section 3.1.2, respectively.

In figure 3.2, a typical haloscope setup is sketched. When placing the optical system in a magnetic field, the axion field sources the effective current $\mathbf{J}_a$ which in turn induces the fields $(\mathbf{E}_a, \mathbf{H}_a)$. The exact shape of $(\mathbf{E}_a, \mathbf{H}_a)$ is determined by the boundaries of the optical system. For calibration purposes, a reflectivity measurement of the optical system is often performed by connecting a Vector Network Analyzer (VNA). The VNA emits a signal of amplitude E_r^+ and measures the reflected amplitude E_r^- . The reflectivity Γ of the optical system is then defined as the ratio of the amplitudes $\Gamma = \frac{E_r^-}{E_r^+}$. A reflectivity measurement induces the fields $(\mathbf{E}_r, \mathbf{H}_r)$ within the optical system.

Haloscopes attempt to design an optical system that results in a high signal amplitude E_a^- towards the receiver. The resulting signal power can be expressed as the time averaged Poynting vector integrated over the port of the receiver system

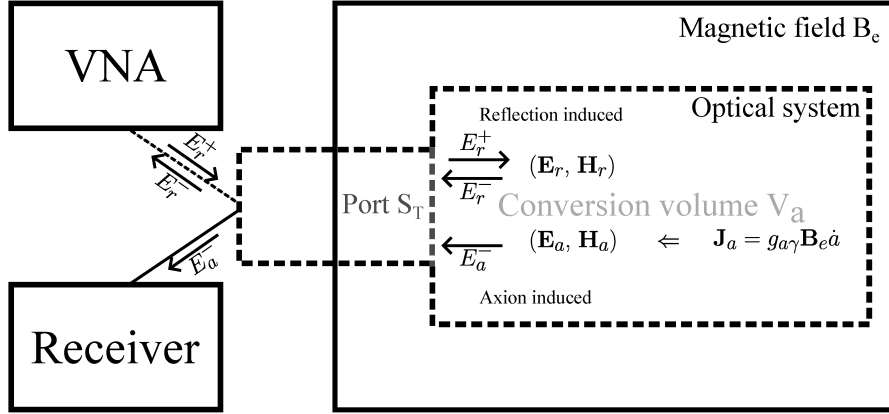


Figure 3.2: Schematic of a typical haloscope experiment, consisting of an optical system and a receiver system. The optical system is placed in a magnetic field $\mathbf{B}_e$, which results in an effective axion current $\mathbf{J}_a$. This current induces the fields $(\mathbf{E}_a, \mathbf{H}_a)$ that potentially result in a detectable signal towards the receiver system E_a^- . Instead of a receiver system, a VNA can be connected to the port of the optical system, which allows to perform reflectivity measurements that induce the reflection-induced fields $(\mathbf{E}_r, \mathbf{H}_r)$.

S_T with the surface normal $\mathbf{n}$ [64]:

$$P_{\text{sig}} = \frac{1}{2} \text{Re} \int_{S_T} dA \mathbf{n} \cdot (\mathbf{E}_a \times \mathbf{H}_a^*). \quad (3.1)$$

Evaluating equation (3.1) requires the knowledge of the fields $(\mathbf{E}_a, \mathbf{H}_a)$ sourced by $\mathbf{J}_a$. In principle, $(\mathbf{E}_a, \mathbf{H}_a)$ can be calculated by solving Maxwell's equations given the boundaries of the optical system but for any real-world system, they will contain unknown geometric imperfections that at least require some measurement-based verification.

To calibrate the system, one would want to experimentally recreate the effective axion current using a local source. However, providing such a uniform current over the extent of the experiment is very challenging. Instead, the sensitivity is typically determined via a reflectivity measurement and the corresponding reflection-induced fields $(\mathbf{E}_r, \mathbf{H}_r)$. Since Γ depends on all the electromagnetic properties of the optical system, it can generally be used to infer characteristics of the reflection-induced electric field $\mathbf{E}_r$. By reciprocity, equation (3.1) can be rewritten in terms of $\mathbf{E}_r$ [64], resulting in a signal power given by

$$P_{\text{sig}} = \frac{1}{16P_{\text{in}}} \left| \int_{V_a} dV \mathbf{E}_r \cdot \mathbf{J}_a \right|^2, \quad (3.2)$$

which only depends on the input power emitted by the VNA P_{in} , the reflection-induced electric field $\mathbf{E}_r$ and the effective axion current $\mathbf{J}_a$ as given by equa-

tion (2.8). From equation (3.2), the expected signal power can be estimated for the different types of haloscopes, as is done below for the dish antenna and cavity haloscopes.

3.1.1 Dish antenna

A dish antenna uses a magnetized mirror to increase the expected signal power from the axion-photon conversion. Over the surface area of the mirror A the magnetic field is assumed to be of constant amplitude $\mathbf{B}_e$ polarized parallel to the mirror. With increasing distance, the amplitude of the magnetic field slowly decreases. Therefore, the integral $|\int dV \mathbf{E}_r \cdot \mathbf{J}_a|^2$ simplifies¹ to $\frac{A}{2}|g_{a\gamma}a_0\mathbf{B}_e|^2$ and equation (3.2) can be written as

$$\begin{aligned} P_0 &= \frac{g_{a\gamma}^2 |a_0|^2 |\mathbf{B}_e|^2 A}{2} \\ &\simeq 8.8 \times 10^{-27} \text{ W} \left(\frac{g_{a\gamma}}{4 \times 10^{-14} \text{ GeV}^{-1}} \right)^2 \left(\frac{|\mathbf{B}_e|}{10 \text{ T}} \right)^2 \\ &\quad \times \left(\frac{100 \text{ } \mu\text{eV}}{m_a} \right)^2 \left(\frac{A}{1 \text{ m}^2} \right) \left(\frac{\rho}{300 \text{ MeV cm}^{-3}} \right), \end{aligned} \quad (3.3)$$

assuming perfect detector efficiency. The values of $g_{a\gamma}$ and m_a correspond to a KSVZ type QCD axion with $C_{a\gamma} = -1.92$ in the mass range motivated by the post-inflationary scenario, as shown in figure 3.1. This signal power is commonly interpreted as photon emissions from the mirror surface.

As will be shown in chapter 4, state of the art receiver systems are sensitive to $\mathcal{O}(10^{-21} \text{ W})$ at room temperature. When cooling the system to cryogenic temperatures, the noise background can be decreased by about two orders of magnitude, resulting in a sensitivity to signals as low as $\mathcal{O}(10^{-23} \text{ W})$. Therefore the signal power needs to be increased further by either stronger magnetic fields, which is technologically challenging, or scaling the surface size, which is quickly limited by mechanical constraints. Still, first limits down to $g_{a\gamma} \lesssim 5 \times 10^{-11} \text{ GeV}^{-1}$ at an axion mass of $m_a \simeq 50 \text{ } \mu\text{eV}$ were already placed [65] using such a detector.

3.1.2 Cavity haloscopes

Cavity haloscopes employ highly resonant microwave cavities that increase the expected signal power from the axion-photon conversion by maximizing the field amplitude of $\mathbf{E}_r$. The resonance of a cavity is commonly characterized using its

¹For an ideal mirror, $\mathbf{E}_r = E_0 \sin(kz)$ with E_0 determined by P_{in} . For any slowly decreasing function $F(z)$, $\int_0^\infty dz \sin(kz)F(z) = \frac{k}{2}$, leaving only a surface integral over some constants.

loaded quality factor Q_L , which for a critically coupled cavity resonant at frequency ν is defined in terms of the stored energy $U = \frac{1}{2} \int_{V_a} dV |\mathbf{E}_r|^2$ as

$$Q_L \equiv \frac{\pi \nu U}{P_{\text{in}}}. \quad (3.4)$$

Furthermore, the normalized geometric overlap between the cavity mode $\mathbf{E}_r$ and the axion current $\mathbf{J}_a$, called the formfactor, needs to be considered:

$$C = \frac{|\int dV \mathbf{E}_r \cdot \mathbf{J}_a|^2}{\int dV |\mathbf{E}_r|^2 \int dV |\mathbf{J}_a|^2}. \quad (3.5)$$

It depends on the operating mode of the cavity. A typical formfactor for a cavity operated under the TM_{010} mode is $C = 0.4$ [58]. The higher the Q -factor, the stronger the resonance, meaning higher sensitivity to the axion. Current cavity experiments typically reach Q -factors in the order of $Q \simeq 10^4$ [66]. The resulting signal power for a cavity of volume V can be written as [64]:

$$\begin{aligned} P_{\text{sig}} &= \frac{g_{a\gamma}^2 \rho}{2m_a} \mathbf{B}_e^2 Q V C \\ &\simeq 3.5 \times 10^{-26} \text{ W} \left(\frac{g_{a\gamma}}{4 \times 10^{-14} \text{ GeV}^{-1}} \right)^2 \left(\frac{\mathbf{B}_e}{10 \text{ T}} \right)^2 \\ &\quad \times \left(\frac{Q_L}{10^4} \right) \left(\frac{V}{10 \text{ cm}^3} \right) \left(\frac{100 \text{ } \mu\text{eV}}{m_a} \right) \left(\frac{\rho}{300 \text{ MeV cm}^{-3}} \right). \end{aligned} \quad (3.6)$$

For simplicity, a critically coupled cavity, resonant at the signal frequency ν with an ideal formfactor of $C = 1$, is considered. Cavity searches have placed limits as low as $g_{a\gamma} \lesssim 10^{-15} \text{ GeV}^{-1}$ for axion masses $\sim 5 \text{ } \mu\text{eV}$ but for higher axion masses they require smaller cavities due to the higher signal frequency, which severely limits their sensitivity. The normalization to the small cavity volume of $V = 10 \text{ cm}^3$ is chosen to reflect this scaling, with the photon wavelength $\lambda \simeq 1 \text{ cm}$ at an axion mass of $m_a = 100 \text{ } \mu\text{eV}$. The corresponding signal power is again too small to detect.

The presented approximations clearly show that the existing haloscope experiments are poorly suited to search for the QCD axion at axion masses around $100 \text{ } \mu\text{eV}$. The next section is therefore dedicated to a new approach called the *dielectric haloscope*, that is designed to overcome this limitation.

3.2 Dielectric haloscopes

Dielectric haloscopes are designed to search for DM axions of masses $\gtrsim 40 \text{ } \mu\text{eV}$, as predicted by the post-inflationary scenario. They combine the concept of a dish

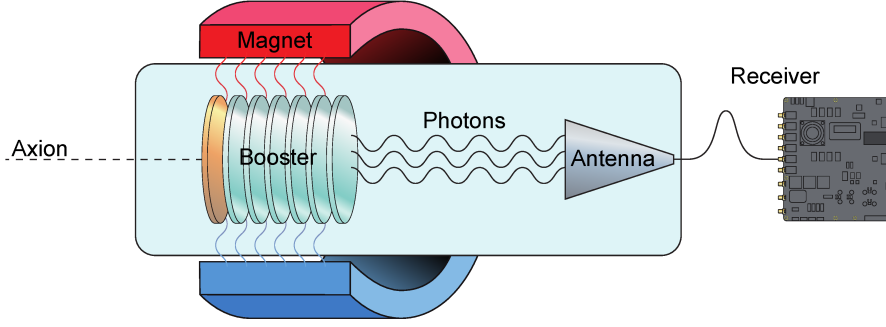


Figure 3.3: Schematic of a dielectric haloscope. It consists of a booster with multiple dielectric disks and a mirror that is placed in a dipole magnet. The signal is coupled to the receiver system using an antenna. Graphic by Olivier Rossel.

antenna with that of a resonant cavity to gain sensitivity to this so far mostly unexplored parameter space. Dielectric haloscopes make use of a *booster* that consists of a stack of dielectric disks and a mirror that enhances the axion-photon conversion within an external magnetic field $\mathbf{B}_e$, as illustrated in figure 3.3. The magnetic field needs to be aligned parallel to the disks and mirror and constant over the extent of the booster, resulting in a uniform effective axion current $\mathbf{J}_a$. P_{sig} (equation (3.2)) can then be maximized by increasing the volume integral of the reflection-induced field $\int dV \mathbf{E}_r$. The resulting signal power is coupled to a receiver system via either an antenna or taper.

The concept is illustrated for a three-disk booster and a one-dimensional field $\mathbf{E}_r$ in figure 3.4. The placement of the three dielectric disks increases the positive contribution of the electric field (red shaded regions) to the signal power by creating a resonance between the mirror and closest disk. The dielectric disks further decrease the negative contributions (blue shaded regions). At frequencies where the setup is not resonant, this still allows for a sensitivity similar to that of a dish antenna of the same radius. By moving the disks, the booster can be tuned for sensitivities at different axion masses.

This signal enhancement can be equivalently described by photon emissions at the interfaces between regions of different permittivity. Given two dielectric media, characterized by their relative permittivities ϵ_1 and ϵ_2 , the axion current would source an electric field in each of those media with an amplitude proportional to $1/\epsilon_{1/2}$. If those two media were now brought in contact, this would lead to a discontinuity of the electric field at the interface. To still satisfy the continuity requirement of Maxwell's equations, photon emissions in opposite directions at the interface are required. Going back to the booster, one can place the disks and mirror in such a way to constructively interfere the coherent emissions from the different interfaces and create resonances. This picture is helpful in one-dimensional

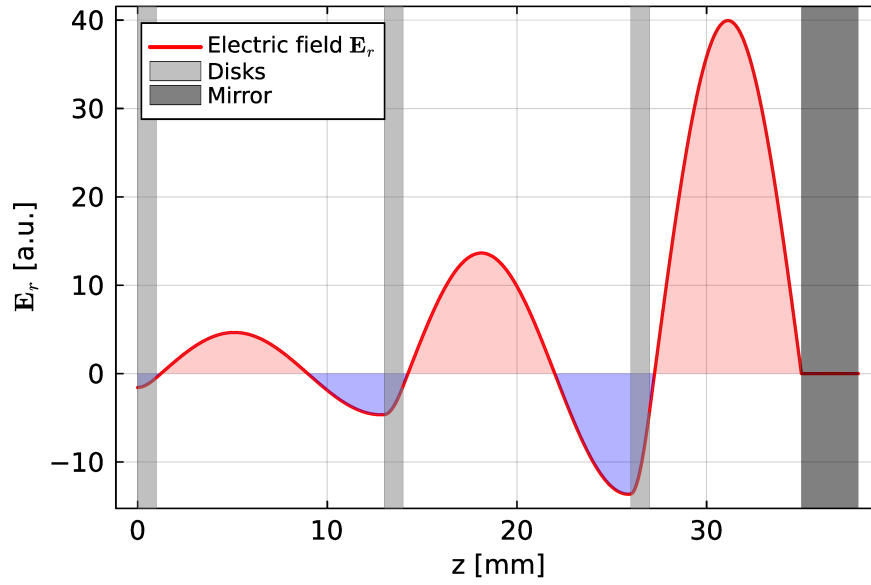


Figure 3.4: Visualization of the reflectivity induced electric field $\mathbf{E}_r$ within a three-disk booster. The disks (light gray) and the mirror (dark gray) provide the boundary conditions that result in the electric field $\mathbf{E}_r$ (red). The red shaded regions indicate positive contributions to the expected signal power P_{sig} and the blue shaded regions negative contributions.

models of the booster. To calculate the expected signal power, the photon emissions can be modeled by plane-waves that are summed up throughout the booster structure, as is done in the model described in section 5.2.1 and then applied to the axion search presented in chapter 6.

To quantify the expected signal power, one defines the boost factor β^2 as the ratio of the expected signal power P_{sig} to the signal power of only a dish antenna with the same surface area as the mirror P_0 (equation (3.3)):

$$\beta^2 = \frac{P_{\text{sig}}}{P_0}. \quad (3.7)$$

Being normalized to P_0 , the boost factor value is independent of booster radius and exact nature of DM - it applies to the dark photon just as much as to the axion. Looking at equation (3.3), a boost factor of $\beta^2 = 10^4$ is required to reach a detectable signal power of $P_{\text{sig}} \simeq 8.8 \times 10^{-23} \text{ W}$. Such a boost factor is expected to be achievable with a booster consisting of twenty dielectric disks [67]. However, it is important to note that equation (3.3) makes some simplifying assumptions, such as assuming a perfect detector efficiency, a disk and mirror surface area of 1 m^2 , which is yet to be built, and requires a magnetic field of $\mathbf{B}_e = 10 \text{ T}$, for which as of now no magnet of sufficient size exists. Nevertheless, this shows the potential of such a setup to reach the desired QCD axion parameter space.

The boost factor is the main measure of sensitivity to DM of a dielectric haloscope. Its determination is however challenging. Due to their large radius compared to the photon wavelength λ of $> \mathcal{O}(10\lambda)$, dielectric haloscopes are highly overmoded systems. Huge care has to be taken to minimize the excitation of higher-order modes (HOMs), because they typically have a much lower overlap with the axion current and therefore lead to a decrease in signal power. This leads to very stringent mechanical requirements [67] because any imperfection results in further HOMs being excited. Furthermore, the existence of HOMs severely impacts the capability to relate reflectivity measurements to the reflection-induced field $\mathbf{E}_r$ inside the booster. Existing methods, like the Q -factor formalism, typically rely on precisely knowing the operating mode of the system. Therefore, new methods are needed for the dielectric haloscope. One of the core contributions of this thesis is the development of a model-based method to determine the boost factor of a closed dielectric haloscope, presented in chapter 5.

With the presence of HOMs, full-wave simulations become less and less reliable, since they need to replicate the exact imperfection that lead to the HOM in the first place. While cavity experiments are operated under a very narrow resonance of a clearly defined eigenmode, one of the defined goals of dielectric haloscopes is to allow a more broadband operation, similar to a dish antenna. This requires to closely study HOMs and eventually quantify their effect on the boost factor, instead of trying to avoid them. Such a method was used in the dark photon

search presented in chapter 7, based on a measurement of the reflection-induced electric field $\mathbf{E}_r$ over the volume of the booster [4].

3.2.1 The Magnetized Disk and Mirror Axion Experiment

The dielectric haloscope concept was originally proposed by the MADMAX working group [23]. It led to the founding of the MADMAX collaboration in 2017 with the goal of reaching sensitivity to the QCD axion within the mass range from $40\text{ }\mu\text{eV}$ to $400\text{ }\mu\text{eV}$. As the baseline design, a booster with 80 disks made from lanthanum aluminate with a surface area of $\sim 1\text{ m}^2$ and a permittivity of $\epsilon = 24$ that is placed in a 9 T dipole magnet was proposed. This is expected to enable boost factors of $\beta^2 \simeq 3 \times 10^4$ over a bandwidth of $\sim 50\text{ MHz}$ [67]. As will be explained in section 4.3, the resulting signal power from a QCD axion, as estimated by equation (3.3), can be detected within $\sim 3\text{ h}$. By moving the disks using piezoelectric motors, the boost factor can be shifted in frequency to cover the planned mass range. To cover the entire range, corresponding to 10 GHz to 100 GHz , approximately $90\text{ GHz}/30\text{ MHz} = 3000$ such measurements have to be performed. This assumes an overlap between boost factors of 10 MHz on each side. Therefore, covering the entire range with the full scale MADMAX setup is possible within a time period of the order of $3000 \cdot 3\text{ h} \simeq 1\text{ yr}$. This assumes continuous data taking with no deadtime, which is clearly not realistic, but the order of magnitude should be a solid estimate.

On its way to reach this goal, several prototype setups have been built, that are categorized as closed booster and open booster setups:

- **Closed booster:** The disks and mirror of closed setups are encased in a metal cylinder, providing fixed boundary conditions that enable a finite number of propagating modes. The signal is then coupled to the receiver system using a taper. Resonances are easy to spot within a reflectivity measurement as distinct dips in reflected power. In a frequency region where only one such feature is apparent, a closed booster is well described as a single moded system, as will be shown in chapter 5. It is however challenging to identify which feature corresponds to which resonating mode. This is a fundamental requirement to determine the resulting signal power extracted from the axion field.
- **Open booster:** Open booster setups consist of free-standing disks and a mirror with the resulting DM signal focused onto an antenna using a focusing mirror. As the goal of MADMAX is to provide tuning capabilities by independently moving the different disks, this is set as the target MADMAX design. It is however challenging to simulate due to its large size and open

boundaries, enabling an infinite set of propagating modes. Furthermore, the antenna reflects any HOM that was excited within the system, leading to a resonant enhancement of those modes. This drastically complicates the modeling of an open booster.

Initially, different open booster setups were used to study the electromagnetic response and mechanical requirements of such systems [25] and design and test the mechanics to move the disks with the necessary $\mathcal{O}(\mu\text{m})$ precision [26]. The closed booster systems were developed specifically to simplify the modeling of the electromagnetic response. This design enabled the boost factor determination using a single-mode model, which is a core contribution of this thesis and described in chapter 5. The successful boost factor determination subsequently enabled a search for axion DM with a closed booster, which is presented in chapter 6. This is an important step towards fully understanding and characterizing more complicated setups. However, the necessary size and tuning capabilities to reach the designated goal of MADMAX make it unsuitable as the final design. Demonstrating the capability to determine the boost factor of an open system is therefore crucial for more sensitive DM searches. Such a method was developed in the PhD thesis by J. Egge [27]. In combination with the noise model developed as part of this thesis, it led to the dark photon search presented in chapter 7.

3.3 Relation to dark photon searches

The theory behind dark photons differs significantly from the axion. However, it results in a coupling to the SM that is experimentally very similar. The equations shown in this section can also be applied to the dark photon by simply replacing $\mathbf{J}_a$ and $(\mathbf{E}_a, \mathbf{H}_a)$ by the corresponding terms for the dark photon $\mathbf{J}_{\text{dp}}$ and $(\mathbf{E}_{\text{dp}}, \mathbf{H}_{\text{dp}})$. This similarity often allows limits placed on the axion parameter space to be converted to a limit on the dark photon parameter space by a simple rescaling. Experimentally, it is sometimes challenging to procure a magnet of sufficient size and strength to search for an axion. In this case, performing a dark photon search instead is a common method to prove experimental viability.

4 Receiver system for a dielectric haloscope

In chapter 3, the signal power emitted by a dielectric haloscope was established. This chapter addresses how this signal power can be detected, which provides the missing piece to the determination of the experiments sensitivity.

Searching for a DM signal from a dielectric haloscope requires a microwave receiver system sensitive to powers of $\mathcal{O}(10^{-21} \text{ W})$ and capable of resolving the narrow spectral features of the DM signal described in section 4.1. The DM-photon conversion is expected to result in a continuous signal at a frequency corresponding to the DM particle's mass. Since the DM mass is not known a priori, the receiver system has to measure spectra of power over frequency. The resolution bandwidth (RBW) should be smaller than the expected signal width $\Delta\nu_{\text{dm}}$. Such a measurement is commonly performed using a spectrum analyzer (SA). However, SAs typically sweep the configured bandwidth in RBW-sized steps resulting in significant deadtime. To avoid this limitation, a heterodyne detection scheme is employed, which is explained in section 4.2.

Crucially, the detectable signal power is limited by the noise introduced by the receiver system. A detection requires a signal-to-noise ratio (SNR) above a certain threshold, commonly set to $\text{SNR} \geq 5$. In section 4.3, the SNR of the experiment is derived, leading to an expression for the sensitivity of a dielectric haloscope to axion dark matter. Furthermore, strategies to reduce the receiver system noise and improve the SNR are discussed.

Unwanted signals present a major challenge, because they can obscure or mimic a potential DM signal. In section 4.4 different sources of unwanted signals are discussed and some methods to suppress them are explained. Their implementation was a key contribution to the dark photon search presented in chapter 7.

Another significant effect arises from reflections at the receiver input, specifically at the amplifier stage. Achieving a low noise figure requires a trade-off with impedance matching [68], which becomes increasingly severe at frequencies above 10 GHz. Such mismatches cause partial reflections of the DM signal and therefore alter the boost factor as discussed in section 4.5. Quantifying this effect constitutes a central contribution of this thesis and forms part of the publication presented in chapter 5. The method is applied to both, the axion search (chapter 6) and the dark photon search (chapter 7) leading to corrections to the boost factor of

$\mathcal{O}(10\%)$.

4.1 Dark matter signal

Haloscopes search for a DM signal from the galactic DM halo. The DM particles are assumed to convert to photons within the optical system. This results in a continuous signal with a frequency ν given by the DM mass m_{dm} and momentum p_{dm} :

$$\nu = \sqrt{p_{\text{dm}}^2 + m_{\text{dm}}^2}. \quad (4.1)$$

The DM particle is assumed to be non-relativistic and therefore the momentum part is typically neglected when presenting plots with one axis corresponding to the DM mass. The momentum is however important when searching for a DM

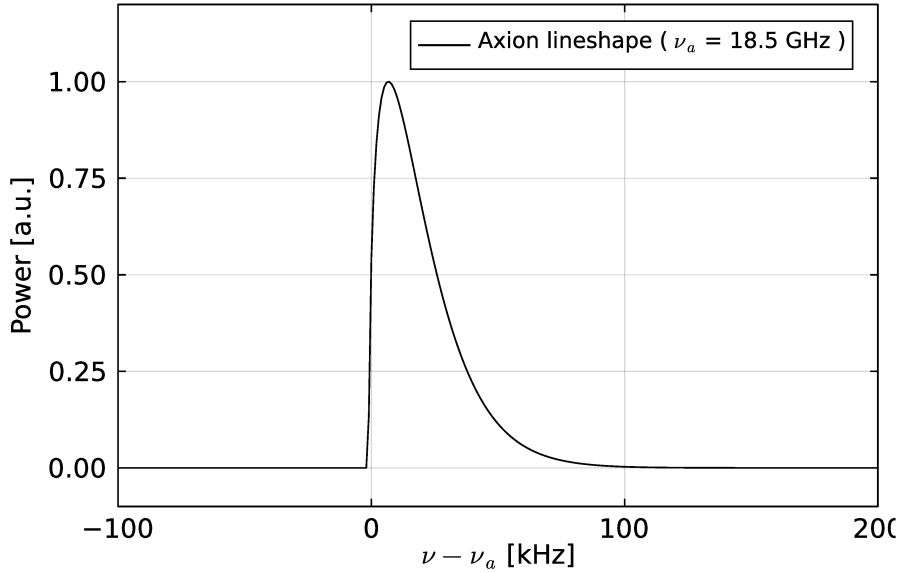


Figure 4.1: Expected lineshape of a DM signal from the galactic DM halo at a frequency of $\nu_{\text{dm}} = 18.5$ GHz, corresponding to a DM mass of $m_{\text{dm}} \simeq 77 \mu\text{eV}$.

signal in a spectrum of power over frequency. While $p_{\text{dm}} = 0$ would result in a δ -peak at ν , in reality the signal is smeared out slightly. Assuming a fully virialized DM halo, the expected velocity distribution of the DM particles is given by a Maxwell-Boltzmann distribution. The resulting approximate width of the DM signal $\Delta\nu_{\text{dm}}$ is then given by the velocity dispersion $\sigma_v \simeq 156 \text{ km s}^{-1}$ [69] as $\Delta\nu_{\text{dm}} = (\sigma_v/c)^2 \nu \simeq 3 \times 10^{-7} \nu$. The distribution further needs to be boosted into the reference frame of the Earth. This yields a distinct signal shape, as shown in

figure 4.1, that can be used for further signal discrimination, given a high enough frequency resolution of the receiver. A more thorough discussion of the signal shape and its derivation is part of the PhD thesis by J. Diehl [70].

4.2 Heterodyne detection

In the presented DM searches, the optical system is connected to an analog to digital converter (ADC), that streams time domain data to a device performing an online fast fourier transformation (FFT) and saving the resulting power spectra. Such a setup enables continuous data takings with close to no deadtime. However, an ADC of a given sampling rate s is only able to digitize signals of a frequency up to its nyquist frequency $\nu_{\text{ny}} = s/2$. Sampling a signal with a frequency > 10 GHz would therefore require an ADC with $s \geq 20 \text{ GS s}^{-1}$, which poses many technical challenges and is as such not a viable option. A heterodyne detection scheme circumvents this by shifting the signal from its original frequency > 10 GHz to a lower frequency range $< \nu_{\text{ny}}$. This is done in hardware using a mixer that multiplies the potential DM signal at frequency ν with a local oscillator (LO) of frequency ν_{LO} . Multiplying two signals of different frequencies results in signals at the frequency difference and sum, based on the trigonometric identity:

$$\cos(2\pi\nu t) \cos(2\pi\nu_{\text{LO}} t) = \frac{1}{2} [\cos(2\pi(\nu - \nu_{\text{LO}}) t) + \cos(2\pi(\nu + \nu_{\text{LO}}) t)]. \quad (4.2)$$

The LO frequency needs to be chosen such that $\nu - \nu_{\text{LO}} \leq \nu_{\text{ny}}$. This low-frequency region is called the intermediate frequency (IF) band, while the frequency region of the signal is called radio-frequency (RF) band. A typical setup is sketched in figure 4.2. Besides the mixer and the ADC, an amplification stage that typically consists of multiple amplifiers to overcome the noise floor of the ADC is required. Using this method, the ADC sampling rate only limits the RF bandwidth to a maximum of ν_{ny} , no longer the signal frequency.

The RBW of such a setup is given by the bin width $\Delta\nu$ of the power spectra. It is determined by the sampling rate s and the buffer size n of the performed FFT as $\Delta\nu = \text{RBW} = \frac{s}{n} = \frac{\nu_{\text{ny}}}{2n}$. It is in general possible to trade bandwidth for a better resolution. By skipping every second data point of the ADC, one can artificially half the sampling rate, which halves the bandwidth but also improves the resolution by a factor two. This concept is known as decimation. It can be useful if the FFT is performed on a device with limited resources such as a field-programmable gate array (FPGA) with a fixed buffer size due to hardware constraints.

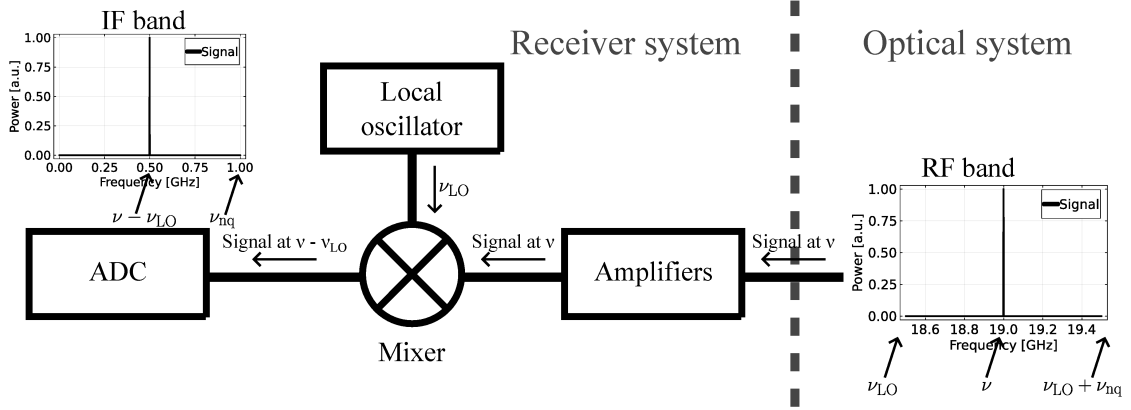


Figure 4.2: A typical setup of a heterodyne receiver system. The optical system on the right produces a signal in the RF band at frequency ν . The signal passes an amplification stage and gets multiplied by the mixer with the signal of the local oscillator at frequency ν_{LO} . Due to the multiplication, the ADC sees a signal in the IF band at $\nu - \nu_{LO}$.

4.3 Signal-to-noise ratio

When searching for any kind of signal, one needs to make sure that the signal-to-noise ratio (SNR) is sufficiently high for detection. The optical system described in chapter 3 is responsible for increasing the signal power, while the receiver system has to minimize the added background noise. At the frequencies of interest > 10 GHz, the main background is expected to be thermal radiation. It is quantified as the frequency dependent system temperature T_{sys} of the system, which is the temperature at which a black body would emit an equivalent noise power. Importantly, in the relevant frequency range $10 \text{ GHz} < \nu < 100 \text{ GHz}$, the power P corresponding to a given system temperature T_{sys} is given by the Rayleigh-Jeans approximation of Planck's law as

$$P = k_B \Delta \nu T_{sys}, \quad (4.3)$$

with the Boltzmann constant k_B .

The system temperature consists of thermal radiation emitted by the booster itself as well as noise emitted by the receiver system. The thermal radiation emitted by the booster is an inherent property of the system and can only be reduced by going to lower physical temperatures, motivating the use of cryostats for DM searches. The receiver system noise is given by the noise emitted by the amplifier stage. If multiple amplifiers with noise temperatures T_i and gains G_i are connected in a chain, the combined noise temperature T_e is given by [71]

$$T_e = T_1 + \frac{T_2}{G_1} + \frac{T_3}{G_1 G_2} + \dots \quad (4.4)$$

The receiver system noise is therefore dominated by the noise of the first-stage amplifier, since all noise sources afterwards are suppressed by its gain. Being an active device, the noise can be reduced by utilizing more sophisticated amplifier devices, referred to as low noise amplifiers (LNAs). Their noise also significantly decreases when going to cryogenic temperatures.

A technologically easier way to improve the SNR is to increase the measurement time. While the DM signal itself is coherent with a coherence time $\tau_{\text{dm}} \propto 1/\Delta\nu_{\text{dm}}$, thermal noise is incoherent white noise. By averaging power spectra over some integration time τ the SNR is expected to increase, as quantified by the radiometer equation [72]

$$\text{SNR} = \frac{P_{\text{sig}}}{k_B T_{\text{sys}}} \sqrt{\frac{\tau}{\Delta\nu}}. \quad (4.5)$$

Equation (4.5) is a central equation for the sensitivity of a dielectric haloscope. As a detection threshold, $\text{SNR} \geq 5$ is commonly required. The noise temperature T_{sys} is expected to be around 400 K at room temperature, consisting of the physical temperature of the booster of 300 K and 100 K of added noise by the receiver system. By replacing P_{sig} with the corresponding signal power determined for a dielectric haloscope in chapter 3 (equation (3.7) and equation (3.3)) and solving for $g_{a\gamma}$, one retrieves the minimal axion-photon coupling that the experiment is sensitive to:

$$g_{a\gamma} \simeq 1 \times 10^{-13} \text{ GeV}^{-1} \left(\frac{T_{\text{sys}}}{400 \text{ K}} \right)^{1/2} \left(\frac{\nu}{20 \text{ GHz}} \right)^{5/4} \left(\frac{10 \text{ d}}{\tau} \right)^{1/4} \left(\frac{\text{SNR}}{5} \right)^{1/2} \\ \times \left(\frac{10 \text{ T}}{\mathbf{B}_e} \right) \left(\frac{1 \text{ m}^2}{A} \right)^{1/2} \left(\frac{300 \text{ MeV cm}^{-3}}{\rho} \right)^{1/2} \left(\frac{3 \times 10^4}{\beta^2} \right)^{1/2} \quad (4.6)$$

The full axion signal power P_{sig} is assumed to enter in one bin, therefore the (configurable) receiver bin width $\Delta\nu$ has been set to $3 \times 10^{-7} \nu$, corresponding to the signal linewidth $\Delta\nu_{\text{dm}}$. The only unknown of equation (4.6) is the boost factor. Its determination for a closed booster is the main topic of the publication included in chapter 5.

Alternatively, equation (4.5) can be used to estimate how long one has to take data to detect a signal of a given power:

$$\tau \simeq 10 \text{ d} \left(\frac{\nu}{20 \text{ GHz}} \right) \left(\frac{\text{SNR}}{5} \right)^2 \left(\frac{T_{\text{sys}}}{400 \text{ K}} \right)^2 \left(\frac{2 \times 10^{-21} \text{ W}}{P_{\text{sig}}} \right)^2. \quad (4.7)$$

When going to cryogenic temperatures, T_{sys} decreases by about two orders of magnitude, therefore, within the same integration time, a signal of $2 \times 10^{-23} \text{ W}$ can be measured instead of $2 \times 10^{-21} \text{ W}$ at room temperature.

However, the integration time τ is not a clear observable of the system: it assumes ideal data acquisition with no deadtime, such that a single power spectrum of resolution $\Delta\nu$ is recorded and averaged every $1/\Delta\nu$. For a more quantitative measure, it is therefore better expressed by the number of averaged power spectra $N = \tau\Delta\nu$:

$$\text{SNR} = \frac{P_{\text{sig}}}{k_B T_{\text{sys}}} \frac{\sqrt{N}}{\Delta\nu}. \quad (4.8)$$

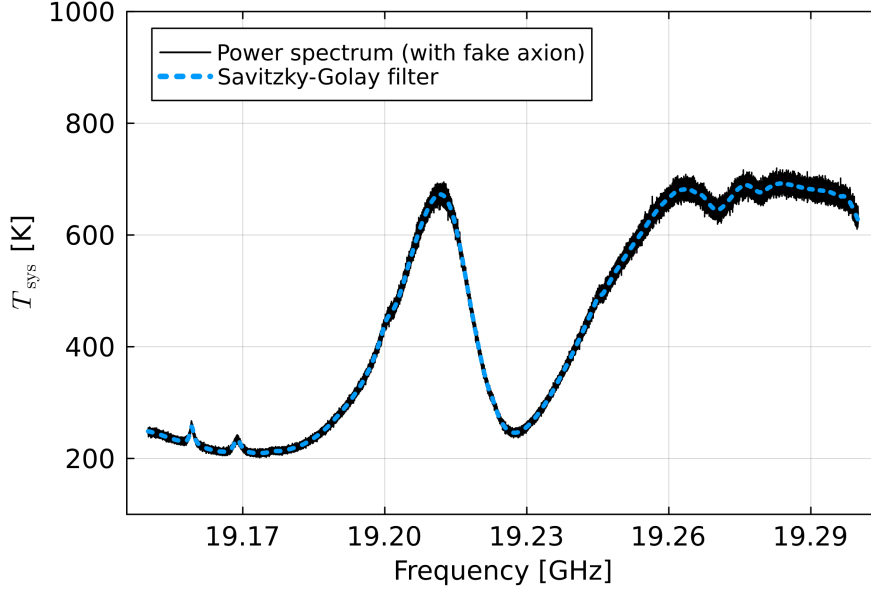


Figure 4.3: Measured power spectrum of a three-disk booster (black solid) with an injected fake axion signal and the same spectrum after applying a Savitzky-Golay filter (blue dashed). For visualization purposes only.

To demonstrate the behavior described by equation (4.8), figure 4.3 displays a power spectrum (black solid) taken with a receiver system, as described in this section, connected to a booster, as described in section 3.2. For illustration purposes, a fake axion is added to the data. It is very challenging to predict the exact background shape a priori, because it depends on the phase behavior of the noise emitted by the receiver system as well as the electromagnetic response of the booster. It is therefore estimated using a Savitzky-Golay filter [73] (blue dotted), which essentially performs a piecewise polynomial fit. The parameters of the filter have to be determined for each combination of receiver and booster individually. This procedure is explained in the thesis by J. Diehl [70]. The presented example is only for visualization purposes. For a sufficiently low polynomial degree and large window size, the filter is not able to resolve the signal lineshape of width $\Delta\nu_{\text{dm}}$ and therefore only fits the background. This is apparent in the normalized

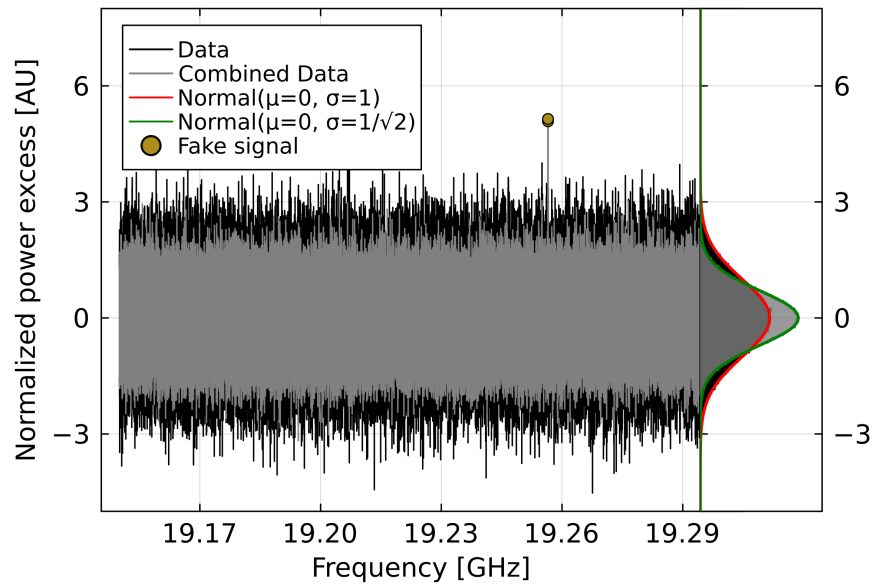


Figure 4.4: Normalized power excess of a single power spectrum (black) and two combined power spectra (gray), including a fake axion signal. The power excess of the single spectrum follows a normal distribution of width $\sigma = 1$ (red) and the power excess of the two combined spectra follows a normal distribution of width $1/\sqrt{2}$ (green). The brown circles indicate the amplitude of the fake axion signal. For visualization purposes only.

power excess, later referred to as processed spectrum p_{proc} , shown in figure 4.4. It is defined as

$$p_{\text{proc}} = \frac{T_{\text{sys}} - T_{\text{SG}}}{T_{\text{SG}}}. \quad (4.9)$$

After background subtraction and normalization (black solid), the injected axion signal is clearly visible, demonstrating that it has not been removed by the filter. The corresponding signal attenuation has to be studied for every setup individually, as is done in section 6.4 and section 7.4 and described in detail in the PhD thesis by J. Diehl [70]. The resulting distribution closely follows a Gaussian of mean $\mu = 0$ and a standard deviation of $\sigma = 1$ (red solid), confirming that the noise is indeed proportional to T_{sys} , as apparent from equation (4.8). When now combining one such normalized spectrum with a second spectrum by performing a bin-by-bin average (gray solid), the combined noise follows a normal distribution with a standard deviation of $\sigma = 1/\sqrt{2}$ (green solid), with the signal amplitude staying about the same. The SNR was therefore increased by a factor of $\sqrt{2}$, confirming the proportionality to $\sqrt{N}$ as predicted by equation (4.8). To ensure proper data taking, the standard deviation of the normalized spectra was monitored continuously for the searches presented in chapter 6 and chapter 7.

4.4 Suppression of unwanted signals

While heterodyne detection is a great tool to digitize signals at high frequencies, some care has to be taken to prevent unwanted signals from entering the data. Such signals can either enter in the RF band, the IF band or be an effect of a non-ideal receiver system.

4.4.1 Signals in the RF band

In the worst case, an external signal (for example from radar systems or similar) that looks similar to an axion or dark photon signal enters in the RF band. There is no general way of avoiding such signals and it always leads to a complete loss of sensitivity in the bins into which the signal enters. One has to verify whether such a signal is a DM signal on a case by case basis. When searching for an axion, this can be as simple as switching off the magnet and checking whether the signal is still visible. For the dark photon, one can verify whether additional shielding around the optical system attenuates the signal – in this case it cannot be a DM signal, which would be generated within the optical system. A slightly more general way is to use a second receiver system during data taking that is only connected to an antenna, without any booster system. If this receiver also picks up a signal,

it is very likely that the signal seen in the actual experiment is also the external, non-DM signal.

4.4.2 Signals in the IF band

Since the receiver system digitizes data from the IF band, signals entering in this frequency range can also be mistaken for a DM signal. The IF band typically covers a big part of the so-called L band, which is in use by many different applications, ranging from mobile service to satellite navigation. Luckily, there is an easy way to distinguish such signals from a potential DM signal in the RF band: while the IF frequency of signals in the RF band depends on the LO frequency, the frequency of signals entering in the IF band does not. To identify unwanted signals from the IF band, one can take two power spectra at slightly offset LO frequencies, each with an integration time of τ_m . For both of them, a list of bins containing a significant signal above 5σ can be generated. Bins that are included in both lists contain IF signals and should be masked for the remaining data run. To which value τ_m should be set will be discussed below.

By combining this with a frequency hopping technique, the masked bins can still be populated with a slightly reduced amount of data. It is based on changing the LO frequency ν_{LO} by some multiple of the FFT bin width $\Delta\nu$ at a fixed time interval τ_{LO} . The number of bins the LO frequency is shifted by n_{shift} is drawn from a normal distribution with a mean of zero and a width of n_{bins} . This method avoids introducing correlations between neighboring bins. Changing the LO frequency shifts signals entering in the RF band while signals in the IF band stay in the same bin. This is the opposite of what is wanted. To inverse the effect, the bins of every recorded power spectra are subsequently shifted by $-n_{\text{shift}}$. To be able to perform a bin-by-bin average of the recorded power spectra, they need to be of the same frequency range, therefore the bins at the edges are rolled over, keeping the number of bins constant. During the analysis the first and the last $\max(n_{\text{shift}})$ bins can therefore not be used. In practice, this is not an issue since the analysis range is given by the narrowest bandpass filter, which usually has a passband sufficiently far from the edges of the spectrum.

To set n_{bins} , τ_m and τ_{LO} to sensible values, it is useful to estimate the decrease of significance of unwanted signals and the introduced reduction of statistics due to the bin masking.

The reduction of statistics can be estimated by the amount of samples taken at the center frequency of a masked bin, which is approximately

$$N_{\min} = N - c \cdot \frac{N}{\sqrt{2\pi n_{\text{bins}}}}, \quad (4.10)$$

with the total number of recorded spectra N and c an integer constant to ac-

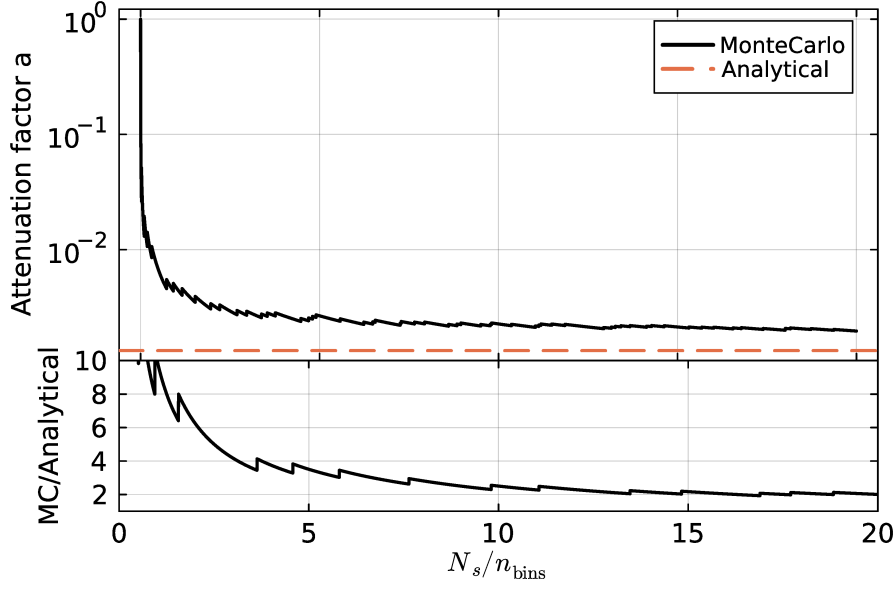


Figure 4.5: MonteCarlo simulation (black solid) and analytical prediction (orange dashed) of the attenuation factor due to the frequency hopping of a signal entering in the IF band. The bottom panel shows the ratio of the two quantities. The match between MonteCarlo simulation and analytical prediction improves with increasing integration time, corresponding to a larger number of samples N_s .

commodate for multiple neighboring bins being masked due to either multiple or broader signals. Since the SNR scales with $\sqrt{N}$, it is estimated to be reduced by a factor of

$$\sqrt{\frac{N_{\min}}{N}} = \sqrt{1 - \frac{10}{\sqrt{2\pi}n_{\text{bins}}}}. \quad (4.11)$$

where $c = 10$ was used, corresponding to a typical number of masked bins observed within the width of the distribution of LO frequencies. This results in a negligible reduction in sensitivity of less than 1 % for $n_{\text{bins}} \gtrsim 300$.

Given sufficient statistics, the attenuation factor a of an IF signal is approximately given by the probability that the LO frequency will be at the mean of its normal distribution as

$$a = \frac{1}{\sqrt{2\pi}n_{\text{bins}}}. \quad (4.12)$$

To estimate what constitutes sufficient statistics, figure 4.5 shows a comparison with a MonteCarlo simulation of a signal with $\text{SNR} = 5$ in a single spectrum after performing $N_s = \tau/\tau_{\text{LO}}$ frequency hopping steps with $n_{\text{bins}} = 300$. It shows that $N_s \simeq 5n_{\text{bins}}$ leads to a deviation by a factor of ~ 4 , while $N_s > 10n_{\text{bins}}$ leads to a deviation by a factor of ~ 2 . The noise level in respect to the noise of a spectrum

used for the bin masking decreases with the total integration time τ as $\sqrt{\tau_m/\tau}$. For eventual unmasked signals to stay under the detection threshold, $a < \sqrt{\tau_m/\tau}$ needs to be true. In terms of τ_m , this yields

$$\tau_m > \frac{4\tau}{2\pi n_{\text{bins}}^2}, \quad (4.13)$$

where $N_s > 10n_{\text{bins}}$ was assumed such that the analytical a was only increased by a factor of 2, as suggested by the MonteCarlo simulation. For $n_{\text{bins}} = 300$, $\tau_m > 6$ s is therefore necessary such that masked signals do never go above the given SNR threshold for a data taking time of $\tau = 10$ d.

To summarize the estimations, one should set $n_{\text{bins}} > 300$ to reach a reduction in sensitivity by less than 1 %. The other parameters depend on the estimated total integration time τ . The interval at which the LO frequency is changed τ_{LO} should be chosen such that $\tau_{\text{LO}} < \tau/(10n_{\text{bins}})$ to reach a good attenuation of unwanted IF signals. The integration time of the power spectra used for the bin masking should be set to $\tau_m > 2\tau/(\pi n_{\text{bins}}^2)$ as per equation (4.13).

4.4.3 Signals due to a non-ideal receiver system

For a typical heterodyne receiver, the following sources of unwanted signals need to be considered:

1. **Negative image frequencies:** when sampling the data, the ADC can not distinguish between negative and positive frequencies - this means that after mixing, signals originating from the frequency region $\nu_{\text{LO}} - \nu_{\text{nq}}$ to ν_{LO} are indistinguishable for the ADC from signals within the frequency region ν_{LO} to $\nu_{\text{LO}} + \nu_{\text{nq}}$.
2. **Mixer sidebands:** Multiplying signals in hardware is difficult. One common way is to rely on a diode, whose current I_d is proportional to the exponential of the input voltage V_i : $I_d \propto e^{V_i}$. When feeding the voltage sum of the potential signal and LO $V_i = V_s + V_{\text{LO}}$ to the diode, the resulting current can be expanded to $I_d \propto V_i + \frac{V_i^2}{2} + \frac{V_i^3}{6} + \dots$. The second-order term contains the wanted multiplication $V_s V_{\text{LO}}$, however higher order terms, although quickly suppressed, can result in additional signals at unwanted frequencies, such as $\nu - 2\nu_{\text{LO}}$. Those are called mixer sidebands. Therefore not only signals around ν_{LO} get shifted to the IF band, but also signals around $2\nu_{\text{LO}}$.
3. **Aliasing effects:** Digitizing an analog signal by sampling at discrete points in time requires a sufficiently high sampling rate to resolve the frequency of the signal. When sampling a signal of frequency $\nu > \nu_{\text{nq}}$, the digitized samples are equivalent to a signal at $2\nu_{\text{nq}} - \nu$.

Those issues can be prevented by adding a band pass filter with steep fall offs around the region of interest and a bandwidth no bigger than ν_{nq} in front of the mixer. It makes sure that:

1. No signal below ν_{LO} can enter the data.
2. No signal outside the region of interest can enter the data, specifically from the mixer sidebands.
3. No signal above ν_{nq} reaches the ADC.

Furthermore, an additional low pass filter with a cut-off frequency $< \nu_{\text{nq}}$ is typically added directly in front of the ADC to make sure that also in the IF band no signal above ν_{nq} enters the data. The negative image frequencies can alternatively be avoided by using so-called IQ mixing, that allows to not only measure the amplitude but also the phase of a signal.

4.5 Amplifier mismatch

At the frequencies > 10 GHz that dielectric haloscopes operate at, the first-stage LNA tends to not be very well matched to the $50\,\Omega$ coaxial connections, with devices reflecting $\sim 20\%$ in magnitude being commonly encountered. This directly affects the boost factor. Since the axion signal is coherent, quantifying this effect requires interference effects to be taken into account. The reflection therefore does not directly translate to a signal loss but might even enhance the signal. Quantifying this effect requires the knowledge of the electrical distance between receiver system and booster, which is not directly measurable. Some experiments counter this issue by introducing a circulator between the LNA and the booster. Such a device allows signal flow only in one direction and therefore absorbs the reflection of the LNA and removes the interference effects. It is however incompatible with magnetic fields and therefore has to be placed at some distance from the booster system. This becomes increasingly problematic at higher frequencies, because the loss of the required cable becomes higher, which directly decreases the sensitivity of the experiment. Operating without such a device, with the LNA connected directly to the antenna or taper, is therefore to be preferred. In section 5.2.2 and section 5.2.3, it is described how this effect can be quantified using a noise model of the receiver system and booster. The model is then used in the axion and dark photon searches, chapter 6 and chapter 7, to determine the effect of the receiver system on the boost factor.

5 Sensitivity of a closed dielectric haloscope to axion dark matter

This section is published as "*Sensitivity of a closed dielectric haloscope to axion dark matter*" [1]. In section 5.7, additional unpublished material is presented.

Significance: As has been established, determining the sensitivity of dielectric haloscopes is particularly difficult. One can neither rely on full-wave simulations nor on an axion-like calibration source. Furthermore, the receiver system directly affects the electromagnetic response of the booster. This publication solves the issue of time-consuming simulations for closed booster systems and presents a model to quantify the effect of the mismatched receiver system on the resulting signal power. The presented models have been developed by the thesis author together with the procedure to determine the boost factor and its uncertainties.

5.1 Introduction

In the search for lightweight dark matter (DM), cavity haloscope experiments [74] have been an essential tool for many years now. They operate under the assumption that the DM from the galactic halo converts to photons. Haloscope experiments were originally motivated by the axion, predicted by a solution to the strong CP problem [75, 11, 76]. Nevertheless, they are also sensitive to other non-relativistic, wavelike DM candidates, such as the dark photon (DP).

By including axions in the Standard Model, Maxwell's equations receive an additional current like source term $\mathbf{J}_a = g_{a\gamma}\dot{a}\mathbf{B}_e$ composed of the axion-photon coupling $g_{a\gamma}$, the time derivative of the axion field $\dot{a}$ and the external magnetic field $\mathbf{B}_e$. Similarly, the DP results in an effective current $\mathbf{J}_{dp} = \chi\dot{\mathbf{E}}_\chi$ with the kinetic mixing angle χ and the time derivative of the DP electric field $\dot{\mathbf{E}}_\chi$. The momentum of non-relativistic DM particles can be neglected, leading to a direct relationship between the particle mass m_{dm} and the resulting photon frequency $\nu \simeq m_{dm}c^2/h$. The measurable signal power of a haloscope P_{sig} is given by the overlap between the effective DM current $\mathbf{J}_{dm} = \mathbf{J}_{a/dp}$ and the electric field $\mathbf{E}$ coupled to the detector: $P_{sig} \propto \int dV \mathbf{E} \cdot \mathbf{J}_{dm}$ [64]. Due to the macroscopic de Broglie wavelength

of the considered DM particles λ_{db} of $\mathcal{O}(10 \text{ m})$, $\mathbf{J}_{\text{dm}}$ is approximately uniform over the extent of the experiment. Typically, resonant radio frequency (RF) cavities operated at the fundamental transverse magnetic mode are used to maximize $\mathbf{E}$ and therefore P_{sig} . Because cavity size at the fundamental mode decreases with frequency, higher DM masses lead to smaller volumes and lower signal power, limiting sensitivity. Different approaches exist to overcome this limitation [77, 78], one being the dielectric haloscope [23].

The MAgnetized Disk and Mirror Axion eXperiment (MADMAX) [79, 23, 24] is such a dielectric haloscope designed to access the mass range $40 \mu\text{eV} \lesssim m_{\text{dm}} \lesssim 400 \mu\text{eV}$, corresponding to a photon wavelength of 3 mm to 30 mm. Based on the magnetized mirror axion DM search idea [78], it uses a *booster* composed of a set of parallel dielectric disks and a metallic mirror, as shown in figure 5.1, to enhance a potential axion DM signal. Depending on the specific booster boundaries, the detector couples either to the fundamental transverse-electric or fundamental transverse-electromagnetic mode, polarized parallel to the disks and mirror in direction of the magnetic field. This in principle allows for an arbitrary transverse scaling of disk and mirror diameter to increase the conversion volume. In practice it is however limited by mechanical constraints and machining tolerances. The placement of multiple dielectric disks allows to shape the distribution of the electric field $\mathbf{E}$ in the longitudinal direction, further increasing $\int dV \mathbf{E} \cdot \mathbf{J}_{\text{dm}}$ and therefore P_{sig} . This effectively decouples the conversion volume from the photon wavelength. The resulting signal power P_{sig} can be equivalently described by coherent photon emissions at the interfaces between air and disk regions, allowing for constructive interference and resonances to increase the experiments sensitivity [48]. In relation to the power emitted by a magnetized mirror, P_0 , P_{sig} is enhanced by the frequency dependent boost factor $\beta^2 := P_{\text{sig}}/P_0$. By adjusting the disk and mirror positions its center frequency and width can be tuned.

The sensitivity of a dielectric haloscope to axions is expressed as the axion-photon coupling $g_{a\gamma}$ needed to produce a given signal-to-noise ratio (SNR) within the integration time Δt for the local axion DM density ρ_a [48]:

$$\begin{aligned}
 |g_{a\gamma}| &= 3.5 \times 10^{-11} \text{ GeV}^{-1} \sqrt{\frac{2 \times 10^3}{\beta^2}} \sqrt{\frac{T_{\text{sys}}}{300 \text{ K}}} \\
 &\times \left(\frac{0.1 \text{ m}}{r} \right) \left(\frac{1 \text{ T}}{B_e} \right) \left(\frac{2.2 \text{ days}}{\Delta t} \right)^{1/4} \sqrt{\frac{\text{SNR}}{5}} \\
 &\times \left(\frac{m_a}{80 \mu\text{eV}} \right)^{5/4} \sqrt{\frac{0.3 \text{ GeV/cm}^3}{\rho_a}}.
 \end{aligned} \tag{5.1}$$

Most of the quantities governing $|g_{a\gamma}|$ are either explicitly set by the experimentalist or directly measurable. This includes the target axion mass m_a , the radius

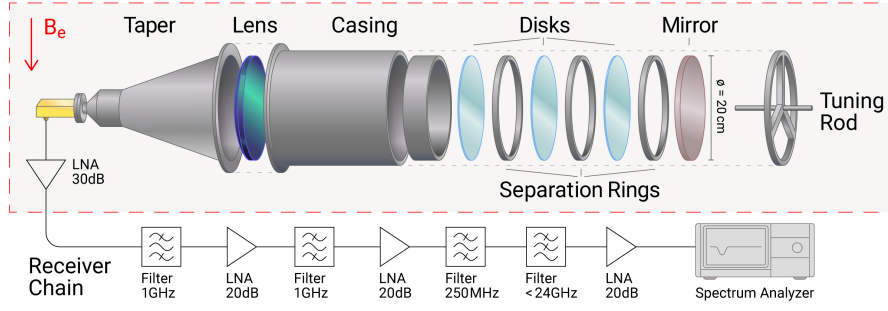


Figure 5.1: Exploded schematic view of the closed dielectric booster CB200 and the receiver chain considered in this paper. In addition to the disks and mirror, placed in an aluminum casing, it includes a taper and a dielectric lens designed to maximize coupling of the axion signal power to the low-noise receiver chain. The tuning rod allows to apply pressure to the mirror, therefore slightly offsetting its position. The shaded region indicates the components exposed to the magnetic field $\mathbf{B}_e$. The setup is identical to the latest axion search performed at CERN’s Morpurgo magnet and the figure taken from its publication [2].

of the disks r and the external magnetic field $\mathbf{B}_e = \hat{y}B_e$, aligned parallel to the disks and mirrors. After performing a standard power calibration [80], described in more detail in section 5.6.1, power spectra taken with the receiver system can be converted to the system temperature T_{sys} . The only remaining parameter is therefore the boost factor β^2 . Its determination requires a more complex procedure since a direct calibration would require to replicate the DM current density $\mathbf{J}_{\text{dm}}$ exactly, which is experimentally challenging. Furthermore, a full-scale simulation of the entire setup is computationally demanding, due to its large domain size of $> \mathcal{O}(10^3 \lambda^3)$ relative to the photon wavelength λ . This stands in contrast to typical cavity domain sizes of $\mathcal{O}(10 \lambda^3)$.

The MADMAX Collaboration has recently published first lower limits on the coupling of axion DM to photons $g_{a\gamma}$ and the kinetic mixing angle χ utilizing two different types of dielectric haloscopes that allow for different approaches to determine β^2 . An *open* booster with three $\varnothing 300$ mm free-standing sapphire disks was used to perform the first MADMAX DP DM search and improved existing limits by nearly three orders of magnitude [3]. In this case, the procedure to determine β^2 was based on a direct measurement of the antenna excited electric field $\mathbf{E}$ within the booster, from which β^2 can be calculated [4]. The measurement of the electric field requires the insertion of a perturbing object into the booster structure and is therefore only applicable to a setup with free standing disks with no surrounding boundary, to which we refer to as an *open* booster. Determining β^2 solely by electric field measurements is very time consuming, which can be improved by relying in parts on a model or simulation.

In a more recent work, a *closed* booster consisting of three $\varnothing 200$ mm sapphire disks encased in an aluminum cylinder, referred to as CB200 and depicted in figure 5.1, marked the first axion DM search [2] with a dielectric haloscope. CB200 uses a dielectric lens and taper to couple the receiver system to the fundamental transverse-electric mode of the booster. The procedure developed for this search only requires access to the port of the taper, relying on a complete model of the booster and connected receiver system, which is the topic of this article. The fixed boundary conditions provided by the aluminum cylinder make it a *closed* booster and ease modeling. As opposed to the open booster, only a finite set of modes propagates within such a system. The search consisted of five periods of data taking, referred to as physics-runs, performed using two different booster configurations in which separation rings of different widths were used. Between physics-runs of the same configuration, the frequency of maximal sensitivity was adjusted by $\mathcal{O}(10 \text{ MHz})$ using the tuning rod, depicted in figure 5.1, to press against the mirror.

This paper presents the analysis that was performed to obtain the already published results [2] in full detail. In section 5.2, a booster and receiver chain model is described and exemplary fitted to the data of the first physics run. With the resulting parameters, the boost factor of the used setup can be calculated. The model is then thoroughly validated against full-wave finite-element-method (FEM) simulations in section 5.3 and applied to the physics-runs in section 5.4, focusing on the resulting uncertainty and time stability of the system. In section 5.5 we provide our conclusions on the performed study.

5.2 Booster and receiver noise model

The frequency dependent boost factor β^2 of a booster is defined by the configuration of the dielectric disks, depending on their positions and permittivities. It is sensitive to $\mathcal{O}(\mu\text{m})$ deviations of not just the disk and mirror positions, but also additional imperfections such as tiny gaps or misalignments. Full-wave simulations for which all of those parameters need to be implemented are very time consuming. Apart from the general computational effort of simulating such large-scale systems, it would also require fine-tuning the parameters describing the imperfections and geometric details to reproduce the exact behavior of the system, multiplying the computational power required for an exact description. Instead, an effective model is used to reproduce the reflectivity of the booster $\Gamma(\nu) = \frac{\phi_{\text{out}}}{\phi_{\text{in}}}$, determined by externally exciting the system by a known field amplitude ϕ_{in} and measuring the reflected field amplitude ϕ_{out} using a Vector Network Analyzer (VNA). The model has to be simple enough to computationally allow for fitting of its multiple parameters. In section 5.2.1 such a model based on network theory [80] is presented.

By fitting its parameters it is able to reproduce the measured reflectivity of the booster $\Gamma(\nu)$ and can then be used to calculate the corresponding $\beta^2(\nu)$. The accuracy of this calculation is discussed in section 5.3.

When taking data for the DM search, the booster is connected to the receiver system via a $50\ \Omega$ coaxial connection. Because the receiver system is not perfectly matched to that connection, it reflects a small fraction of the potential signal, altering the boost factor. With the receiver connected, it is only possible to measure system temperature spectra, consisting of the noise emitted by the receiver and booster. A noise model reproducing such a spectrum is therefore necessary to quantify the effect of a mismatched receiver system on the boost factor. The noise model of the receiver is described in section 5.2.2, requiring additional system temperature measurements of known standards to determine its parameters. It is then combined with the booster model in section 5.2.3 to reproduce the measured system temperature of the booster by fitting the only free parameter left: the electrical length that connects the receiver to the booster.

In the following, the full procedure to determine the sensitivity of a given CB200 configuration is demonstrated for the first physics-run published in [2], referred to as run 1.1. It consists of determining:

1. booster model parameters by fitting to a reflectivity measurement of the booster,
2. receiver noise model parameters by fitting to system temperature measurements of known standards,
3. the electrical distance connecting both models by fitting to a system temperature measurement of the booster connected to the receiver system.

5.2.1 Booster model

We first focus on modeling the booster's reflectivity to extract the parameters necessary to calculate β^2 from a reflectivity measurement. Due to its conducting cylindrical boundaries, wave propagation within the closed booster is well described by the modes of a cylindrical waveguide. In the presence of a homogeneous magnetic field, the axion-induced effective current $\mathbf{J}_a$ excites mainly the fundamental transverse electric (TE_{11}) mode, which propagates along the booster axis, orthogonal to the mirror surface. We call the resonance of this mode the *booster mode*. Such guided wave propagation is modeled using a well-known transmission line (TL) approach [80], implemented in Advanced Design System (ADS) [81]. It is equivalent to using a transfer matrix based method as described in [48], but enables easy integration of the booster into the noise model, also implemented in ADS and described in section 5.2.2 and section 5.2.3.

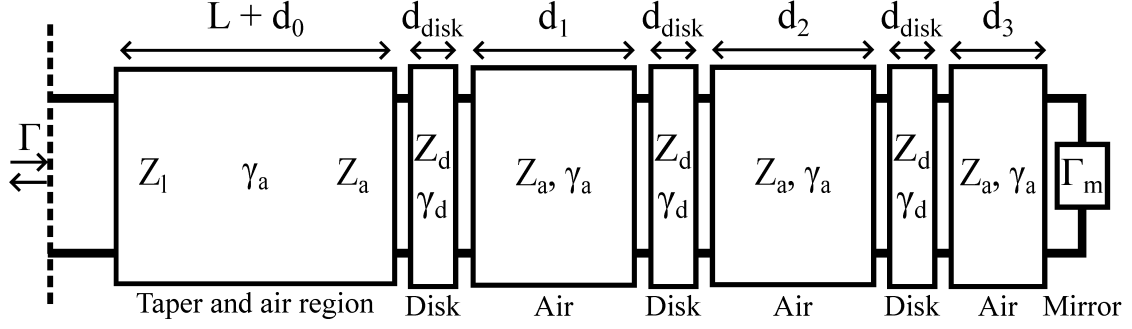


Figure 5.2: Transmission line configuration used to model the booster. It consists of four air and three disk regions of respective lengths d_i and d_{disk} , defined by their corresponding impedances Z_a and Z_d as well as propagation constants γ_a and γ_d . The leftmost air region also includes the length of the taper L . It transitions from Z_a to the coaxial line impedance $Z_l = 50 \Omega$. The mirror is modeled by its reflectivity Γ_m given by its finite conductivity. The reflectivity measurement is modeled at the reference plane indicated by the dashed line, which corresponds to the connection of the VNA to the taper, the booster input.

The model does not consider higher-order mode contributions to the boost factor. It is therefore a fundamental requirement to identify a frequency range in which only the booster mode is present. This requirement is studied in section 5.3.3 and section 5.3.4, using simulations and complementary measurements.

Figure 5.2 shows the booster model, composed of TL segments for the disk (d) and air (a) regions, characterized by their impedances Z_d and Z_a , respectively, their propagation constants γ_d and γ_a and their individual lengths d_{disk} and d_i ($i = 0, 1, 2, 3$). The TL segments directly correspond to the booster parts as shown in figure 6.1. While the disks are all of 1 mm thickness, the separation rings are chosen individually for the different configurations. The electric field of the resonating booster mode is mostly concentrated between the mirror and the closest disk, such that it is most sensitive to the distance between the two, d_3 . The leftmost region includes the measured length of the taper $L = 166(1) \text{ mm}$ and the width of the leftmost separation ring $d_0 = 37.40(1) \text{ mm}$ for configuration 1 and $d_0 = 21.14(1) \text{ mm}$ for configuration 2. The region's length does not affect the boost factor and is therefore kept fixed for each configuration. The taper is assumed to provide a perfect coupling between its coaxial port and the booster mode, implemented by an ideal taper that transitions from the coaxial line impedance $Z_l = 50 \Omega$ to the air impedance Z_a without introducing reflections. This assumption is verified in section 5.3.3.

The equations in this paragraph are based on the description of a circular waveguide in [80], unless stated otherwise. The impedance as a function of the relative

permittivity ϵ in the case of the circular waveguide is given by:

$$Z(\epsilon) = Z_0 \sqrt{1/\epsilon} \frac{k(\epsilon)}{\gamma(\epsilon)}, \quad (5.2)$$

with the free space impedance $Z_0 \simeq 377 \Omega$, dispersion relation $k(\epsilon) = \omega \sqrt{\epsilon}/c$, angular frequency $\omega = 2\pi\nu$ and speed of light in vacuum c . The disk and air impedances $Z_d = Z(\epsilon_d)$ and $Z_a = Z(\epsilon_a)$ differ between air and disk regions only due to their different relative permittivities: $\epsilon_d = 9.36$ for sapphire and $\epsilon_a = 1$ for air. The change in phase per unit length along the booster axis is defined by the propagation constant of the TE_{11} mode:

$$\gamma(\epsilon) = \sqrt{\frac{\omega^2 \epsilon}{c^2} - \frac{p'_{11}{}^2}{r^2}} \quad (5.3)$$

with waveguide radius r , and the first root of the derivative of the Bessel function of order one $p'_{11} = 1.841$.

Loss due to the finite conductivity $\sigma_{\text{Al}} = 3.77 \times 10^7 \text{ S m}^{-1}$ for the aluminum booster walls is described by the attenuation per unit distance:

$$\alpha_c = \frac{1}{rkZ\gamma} \sqrt{\frac{\omega\mu_0}{2\sigma_{\text{Al}}}} \left(\frac{p'_{11}{}^2}{r^2} + \frac{k^2}{p'_{11}{}^2 - 1} \right). \quad (5.4)$$

Additional dielectric loss in the disk segments is accounted for by:

$$\alpha_d = \frac{\tan \delta}{2} \frac{k(\epsilon_d)^2}{\gamma(\epsilon_d)}, \quad (5.5)$$

where a uniformly filled lossy dielectric with loss tangent $\tan \delta$ is considered. Dielectric loss within air regions is neglected. The non-perfect reflection Γ_m at the aluminum mirror due to its finite conductivity is given by $\Gamma_m = 1 - 2\sqrt{2\omega\epsilon_0/\sigma_{\text{Al}}}$ [82]. By connecting transmission line segments as shown in figure 5.2, with propagation constants γ and attenuation $\alpha_c + \alpha_d$ for disks and α_c for air within ADS, an ADS S-parameter simulation is used to calculate the resulting sum of reflections at the interfaces and therefore the reflectivity of the booster $\Gamma(\nu) = S_{11}$.

By fitting the disk spacings d_i , as well as disk permittivity ϵ_d , thickness d_{disk} and loss $\tan \delta$ of the described model, the booster mode visible in the measured reflectivity of the CB200 booster is reproduced within uncertainties, as can be seen in figure 5.3. The model and measurement agree for both the reflected power $|\Gamma|^2$ (top curves) and group delay $\tau_{\text{gd}} = -\frac{\partial \arg \Gamma}{\partial \omega}$ (bottom curves). Since the refractive index of the disks $n = \sqrt{\epsilon_d}$ and the disk thickness d_{disk} are almost fully correlated, the product nd_{disk} , referred to as disk phase thickness, is fitted instead of each

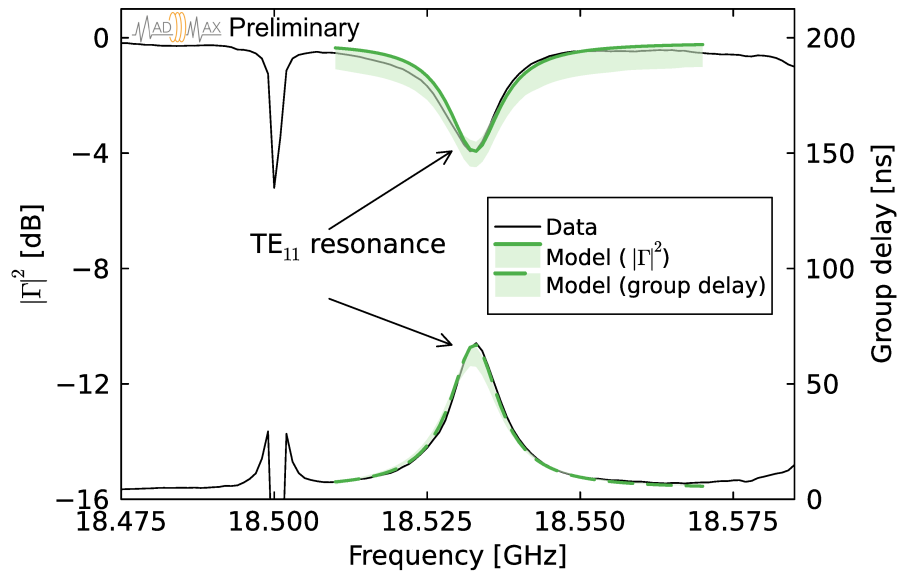


Figure 5.3: Measured (black) and modeled (green) reflectivity of CB200, run 1.1 (see table 5.1), after parameter fit. Top curves: reflected power $|\Gamma|^2$. Bottom curves: group delay τ_{gd} . The model curve shows the best fit and the band the 1σ fit uncertainty within the frequency range of the booster mode. The TE_{11} resonance, that is the booster mode, is visible as a dip in reflected power and a peak in group delay. Fit parameters are listed in table 5.1 and fixed parameters in table 5.2 in the section 5.6.2 .

separately. As initial parameters for the fit, the physical dimensions of the disks and separation rings are used as well as the expected permittivity $\epsilon_d = 9.36$ and dielectric loss $\tan \delta = 10^{-5}$ of sapphire. The fitted parameters, together with their uncertainties, will be discussed in section 5.4. It also contains table 5.1, which includes the parameters used in the fit exemplary presented here for run 1.1. All fixed parameters are summarized in table 5.2 in the section 5.6.2.

To calculate the frequency dependent boost factor, the axion-induced excitation is modeled by including voltage sources of power $P = P_0(1 - 1/\epsilon_d)$ with opposite phase to each side of every disk and one of $P = P_0$ to the mirror, with P_0 being the power emitted by a magnetized mirror [78]. The signal power P_{sig} is then given by an ADS power probe at the input of the booster, resulting in $\beta_{\text{ID}}^2 = P_{\text{sig}}/P_0$. This quantity does not yet take into account the transverse field shape within the booster. While the effective current $\mathbf{J}_a$ due to the axion within the magnetic field is approximately constant over the whole volume, the tangential part of the electric field of the booster mode $\mathbf{E}_{\text{TE}_{11}}$ vanishes at the metallic boundaries.

The coupling is therefore reduced by the transverse geometric overlap of the two, resulting in $\beta^2 = |\eta_A|^2 \beta_{\text{ID}}^2$ with

$$|\eta_A|^2 = \frac{\left| \int_A dA \mathbf{E}_{\text{TE}_{11}} \cdot \mathbf{J}_a \right|^2}{\int_A dA |\mathbf{E}_{\text{TE}_{11}}|^2 \int_A dA |\mathbf{J}_a|^2} \Big|_{\mathbf{J}_a = \text{const}} \quad (5.6)$$

$$= \frac{\left| \int_A dA \mathbf{E}_{\text{TE}_{11}} \cdot \hat{y} \right|^2}{A \int_A dA |\mathbf{E}_{\text{TE}_{11}}|^2} = 0.84. \quad (5.7)$$

This quantity is conceptually similar to the form factor of a cavity, with the difference that the form factor describes the overlap of the resonating mode with the axion current over the volume of the cavity while $|\eta_A|^2$ only considers the transverse overlap, since the longitudinal behavior is already described by the booster model.

5.2.2 Receiver noise model

To take into account the effects of the receiver system on the boost factor, we need to model the system temperature of the combined system of booster and receiver. An ADS noise simulation is used for this purpose. In our case, the dominant noise sources are the first-stage low noise amplifier (LNA) and the booster itself. While ADS automatically simulates noise of passive devices such as the booster according to their signal attenuation, active devices such as the LNA require the addition of specific noise sources.

We model the amplifier noise with a correlated current and voltage noise source, I_n and V_n , connected in parallel [83], as depicted in figure 5.4. The circuit is

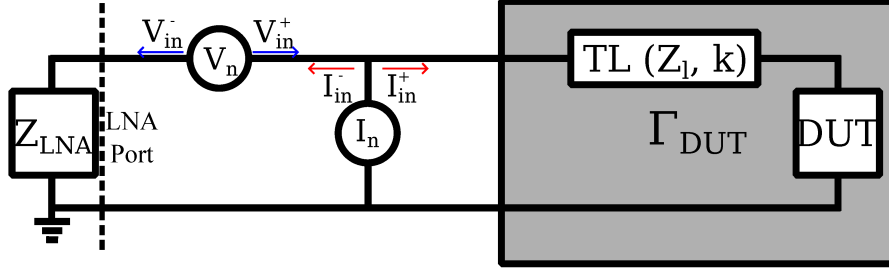


Figure 5.4: Schematic of the circuit used to model LNA noise. It consists of a transmission line of impedance $Z_l = 50 \Omega$ terminated by the LNA impedance Z_{LNA} on one side and a device-under-test (DUT) such as the booster on the other side. The LNA noise is introduced by a voltage and current noise source connected in parallel. The dashed line represents the LNA input port.

terminated by the first stage LNA, implemented as a one-port device with the impedance set to the frequency-dependent impedance of the input port of the LNA Z_{LNA} , measured by a VNA. The device-under-test (DUT) can be described either by its measured or modeled reflectivity Γ_{DUT} . The system temperature as measured with the receiver system is then modeled using the simulated noise power delivered to the one-port device, representing the LNA. Due to the potential correlation of V_n and I_n , this results in four parameters for the receiver noise model: voltage noise amplitude V_n , current noise amplitude I_n , correlation magnitude $|c|$ and correlation angle ϕ_c . They are extracted from fits to system temperature spectra, that are measured with the LNA connected to a short, open and matched load standard as the DUT. Γ_{DUT} is given by their idealized reflectivities $\Gamma_{\text{short}} = -1$, $\Gamma_{\text{open}} = 1$ and $\Gamma_{\text{load}} = 0$. For the open and short standards, the distance between LNA and standard is added as an additional fit parameter. The parameters are fitted to all three spectra simultaneously, with the result shown in figure 5.5 and the fitted parameters listed in table 5.3 in the section 5.6.2.

The noise model reproduces the measurements within their uncertainties, stemming from the receiver power calibration (section 5.6.1). It matches nicely the slightly skewed shape of the oscillation present in open and short measurements. The maximum of the spectrum of the short standard, at around 17.5 GHz, is slightly overestimated and the maximum of the spectrum of the open standard, at around 19.4 GHz, is slightly underestimated. This hints at an increase in noise emitted by the LNA with frequency, but is well within the uncertainties that will be discussed in section 5.4.

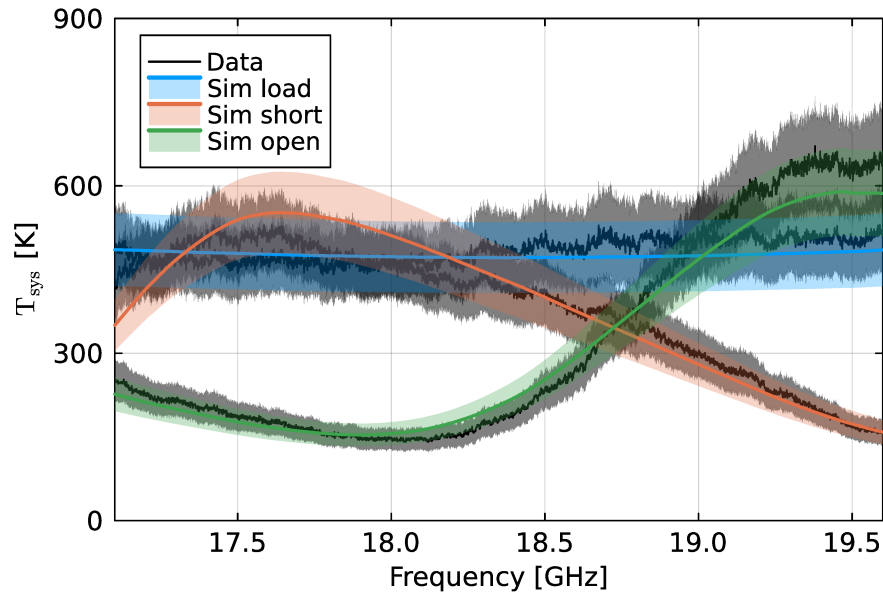


Figure 5.5: Measured and simulated noise spectra T_{sys} of the amplifier connected to open, short and matched load standards. The bands around the curves show their associated uncertainty, stemming from the power calibration of the data. The load (blue) is close to constant over the frequency range, while the open (green) and short (orange) measurements exhibit an oscillation of opposite phase. The simulation matches the measurements within uncertainties.

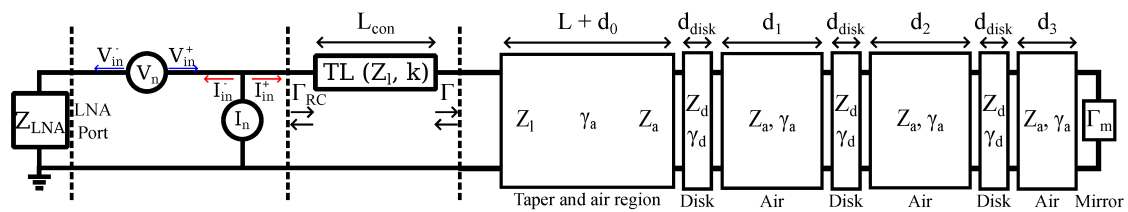


Figure 5.6: Combined noise model consisting of the receiver noise model and booster model connected by a transmission line of length L_{con} . The dashed lines represent the reference planes from which the booster and receiver chain reflectivity measurements, Γ and Γ_{RC} respectively, are taken with the VNA and the LNA input port.

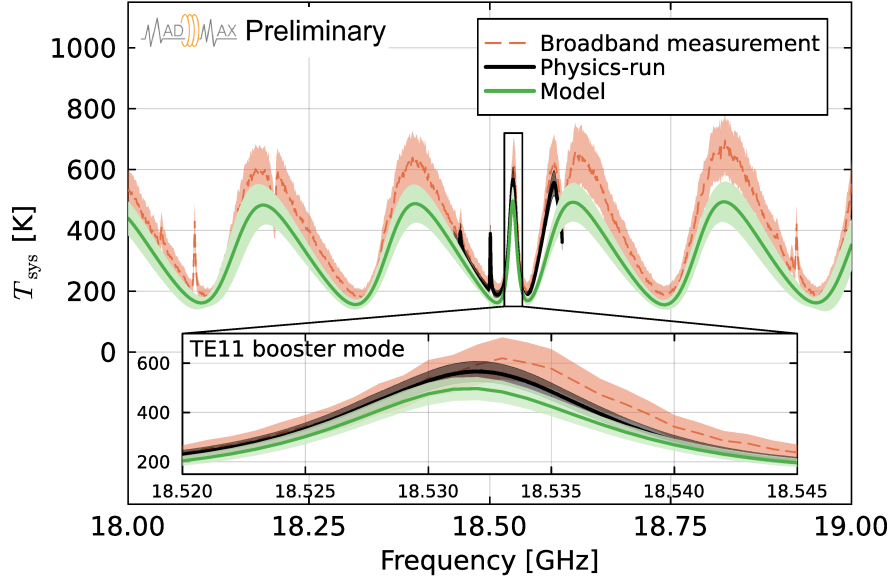


Figure 5.7: Measured and modeled power spectra of CB200, resonant at ~ 18.53 GHz for run 1.1 (see table 5.1). The model (green solid) reproduces the positions of minima and maxima of the broadband oscillation present in the calibration measurement (orange dashed) as well as the booster mode frequency of the physics-run (black solid). The amplitude is consistently slightly underestimated. The inset displays a zoom in on the frequency of the booster mode.

5.2.3 Combined noise model

To combine the LNA and booster model, they are connected by a transmission line of length L_{con} within ADS. This results in the combined model as shown in figure 5.6, where the device under test from figure 5.4 is replaced by the booster model described in section 5.2.1. Due to the impedance mismatch of the LNA to the line, the boost factor as simulated by ADS differs from the simulated one without the LNA. To quantify the change, the only additionally required parameter is the length L_{con} of the transmission line connecting LNA and booster.

Figure 5.7 shows a broadband measurement (1 GHz) of the system temperature of CB200 (orange dashed), the first narrowband spectrum taken during run 1.1 (black) as well as the reproduction by the noise model (green solid).

The broadband measurement features a very prominent oscillation, which is the result of a standing wave between the LNA and the booster. The distance between maxima and minima in frequency space directly depends on the remaining unknown parameter L_{con} . By fitting L_{con} , the positions of minima and maxima over the full 1 GHz range of the measurement are consistently reproduced. Having a broadband measurement increases the precision on L_{con} , since multiple minima

and maxima are visible. This configuration of the receiver system is dedicated to determining the boost factor before the physics-run.

The physics-run spectrum was taken in a narrowband configuration of the receiver system optimized for the axion search. It does not include multiple minima and maxima of the broadband oscillation, but instead allows to spot any changes in booster mode frequency between the time of the broadband measurement and the start of the physics-run. They are visible as a shift in frequency of the feature corresponding to the booster mode, magnified in the inset. In the shown case, the booster mode frequency slightly changed, necessitating an adjustment of the model. As will be further elaborated on in section 5.4.2, the change is caused by an initially unstable mirror position. It is therefore accommodated for by slightly varying the distance between mirror and closest disk d_3 of the model. In this specific case, the required change was only $\sim 1 \mu\text{m}$, corresponding to a $\sim 1 \text{ MHz}$ shift in frequency. Uncertainties will be discussed in section 5.4.

The model does not include higher order modes and therefore does not reproduce additional peaks and dips present in the spectra outside of the booster mode frequency. Furthermore, the amplitudes between broadband measurement and physics-run can change slightly because of the reconfiguration of the receiver system as well as environmental factors such as the temperature. This also affects the amplitude obtained from the model, since it is defined from the previously performed fits to the standards. An overall deviation in amplitude, however, does not affect the boost factor calculation at all, since the parameter extracted from this measurement, L_{con} , only changes the accurately reproduced positions of minima and maxima.

5.3 Model verification

While the presented booster model is able to reproduce the performed reflectivity measurements within uncertainties, it remains to be shown, that the resulting boost factor calculation holds up to scrutiny. We therefore present multiple full-wave FEM simulations to verify that the booster model accurately describes the axion induced excitation and is able to accommodate for 3D effects such as tilts and disk deformation using effective parameters. Those FEM simulations do not include the receiver part of the setup, since it is fully single-mode and therefore completely described by the presented transmission line based receiver model. Previous studies have already been performed on the effect of geometrical inaccuracies of disks on the boost factor of an open booster setup with ≥ 20 disks [67], as opposed to the three disk closed booster considered in this work.

In section 5.3.1 and section 5.3.2 the booster model is compared to those simulations, determining its limits and accuracy. In section 5.3.3 and section 5.3.4 it

is then described how the booster mode is identified and isolated from unwanted modes, which is an inherent requirement for the validity of the booster model.

5.3.1 Ideal case

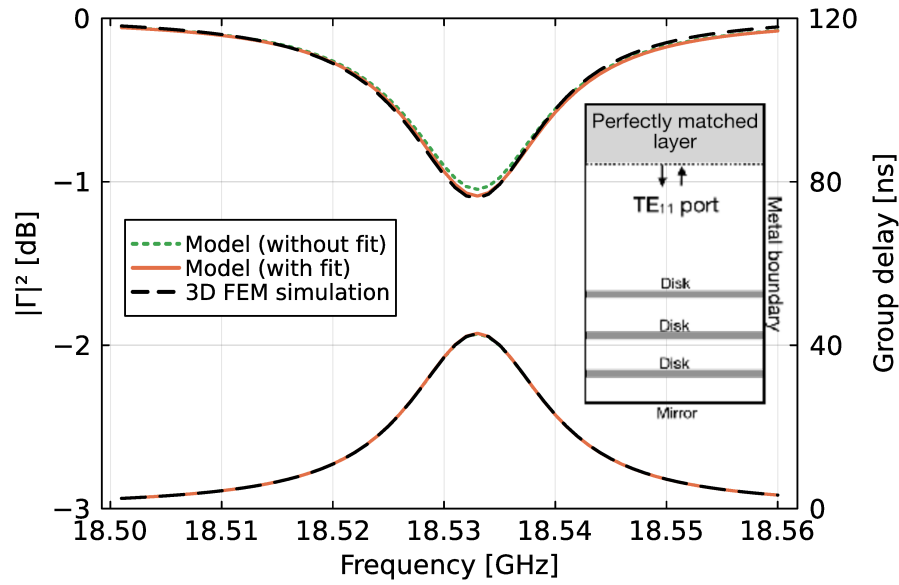
We first consider an ideal booster with perfectly planar disks and a mirror, free of tilts or misalignment, surrounded by a highly conducting boundary and equipped with an ideal, non-reflecting taper supporting only the intended TE_{11} mode. As sketched in the inset of figure 5.8, such a taper is simulated using a perfectly matched layer behind the port. This idealized case is implemented in a full-wave benchmark FEM simulation performed using COMSOL Multiphysics® [84]. A realistic taper, exciting higher-order modes (HOMs), will be discussed in section 5.3.3. Due to computational constraints, the radius of the cylinder is set to 5 cm instead of 10 cm. Since an ideal booster does not excite HOMs, the results are independent of booster radius. The axion-induced excitation is modeled with a background current given by $\mathbf{J}_a$, aligned to the polarization of the TE_{11} mode. The simulated boost factor β^2 is calculated at the TE_{11} port from the power coupled to the TE_{11} mode. External excitation is implemented with a standard TE_{11} port, from which the simulated reflectivity $\Gamma(\nu)$ is obtained.

In figure 5.8 we show the 3D FEM simulation result for an ideal booster with three sapphire disks, that is resonant at ~ 18.55 GHz, similar to the booster in run 1.1 as considered in section 5.2. The comparison to the simulated reflectivity is performed with and without parameter fitting (using exactly the simulated material and geometric parameters).

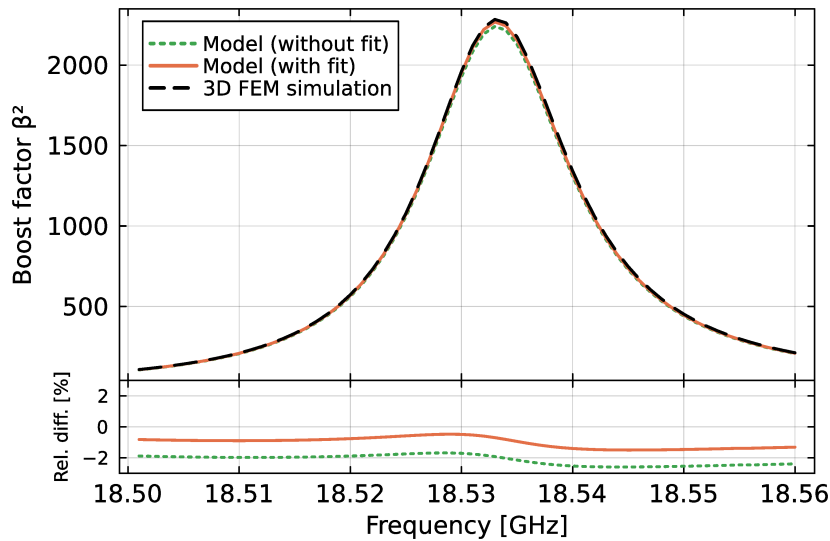
Already without a fit, the group delay obtained with the booster model is indistinguishable from the one obtained with the 3D FEM simulation. Only the reflected power deviates slightly, which is likely caused by a different treatment of conductive losses at the boundaries and the mirror between the booster model and the COMSOL 3D FEM simulation. In terms of boost factor β^2 , an excellent agreement is seen between the model and 3D FEM simulation with $< 3\%$ difference throughout the frequency spectrum without fitting the parameters and $< 2\%$ difference after the fit. These results are consistent across different configurations and demonstrate the capability of the booster model to accurately calculate the boost factor of an ideal booster and already show how the model can make up for slight deviations by adjusting its effective parameters.

5.3.2 Treatment of 3D effects

The booster model used to describe the electromagnetic response of the booster is inherently one-dimensional. Realistic inhomogeneities of material properties, such as the dielectric constant, conductivity or thickness, break the azimuthal



(a) Comparison of reflectivities with schematic of the simulation setup in the inset.



(b) Comparison of boost factors and their relative difference.

Figure 5.8: Comparison of the reflectivity and boost factor obtained from the booster model and the 3D FEM simulation of the setup sketched in the inset of (a) (rotated by 90° for graphical reasons). The spectra of reflected power $|\Gamma|^2$ and group delay τ_{gd} are shown in (a), whereas the boost factor spectra β^2 together with their relative difference are shown in (b). The results obtained from the booster model without fitting its parameters (green dotted) and with fitting its parameters (orange solid) match the results obtained with a 3D FEM simulation (black dashed) of an ideal booster comparable to CB200 during run 1.1 within 3 %.

symmetry of the geometry and redistribute the field. This can lead to a change in boost factor due to: (i) deformation of the intended booster mode and coupling to HOMs; (ii) shift of frequency of the booster mode, also resulting in changes of the reflectivity; and (iii) localized field deviations that alter the coupling to the input antenna. Such three-dimensional effects are accounted for in the booster model through effective parameters. This section demonstrates that the booster model is capable to describe the small discrepancies typically observed between idealized and measured booster responses by allowing the model parameters to deviate from their nominal physical values.

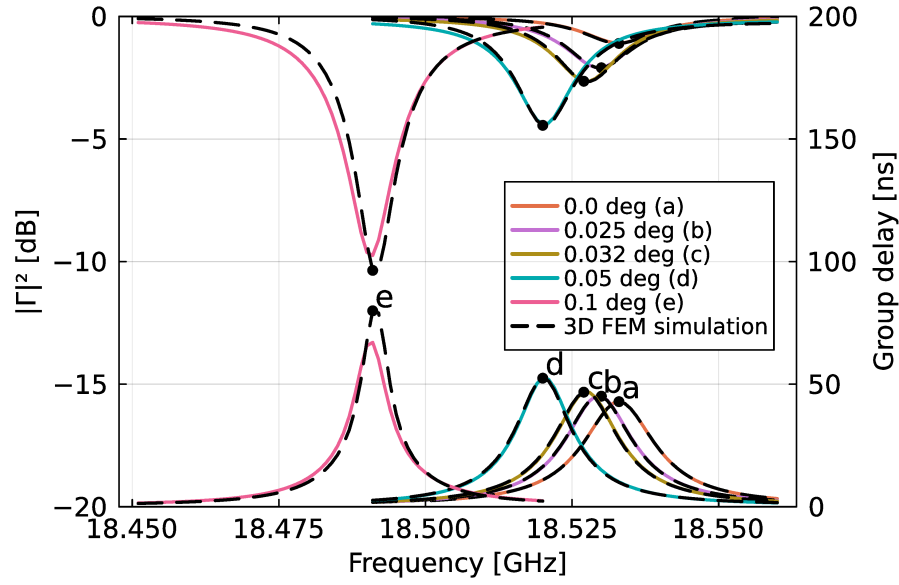
Mirror tilt is the dominant 3D effect on the booster response as the booster is set up for a resonance between mirror and first disk. In practice, the disk spacings are well controlled by precision-machined metal rings, while the mirror position is deliberately left adjustable to enable fine-tuning of the frequency of maximal sensitivity. This design choice provides a useful degree of tunability, but also modifies the booster response. Accordingly the influence of the mirror boundary deserves closer examination.

To realistically account for the effect of the mirror tilt, we modify the 3D FEM simulation described in section 5.3.1 to include a tilted mirror, defined by a rotation around a given axis. The x - and y -axis define the plane of the untilted mirror, with $\mathbf{J}_a$ polarized along the y -axis and the origin at the mirror center. The z -axis is the normal of that plane, pointing towards the taper.

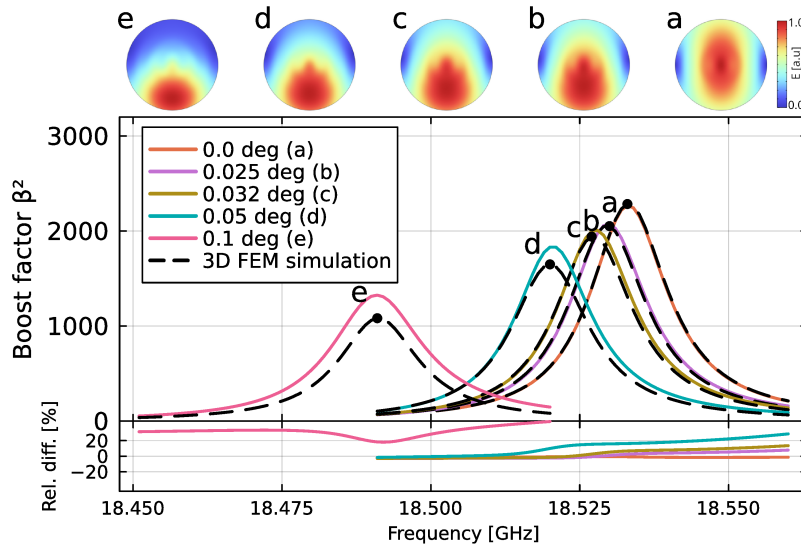
A tilt of the mirror perturbs the field, such that the power in the booster mode reduces due to mixing with unwanted HOMs. The boost factor evaluation only considers the power of the booster mode, excluding power contribution from HOMs. This provides a conservative estimate of the power reaching the detector, since in practice the taper is able to convert some of the HOM power back into the desired mode. The reduced booster radius of 5 cm remains very large compared to the wavelength of ~ 1.6 cm and therefore does not significantly affect this analysis: a radius of 5 cm supports roughly 100 propagating modes at 18.5 GHz, allowing for the field perturbations expected from the small deviations from an ideal booster that we simulate.

Figure 5.9 shows the simulated booster response (black dashed) for a mirror tilt between 0 deg (case a) and 0.1 deg (case e) around the x -axis.

As expected, increasing tilt produces an observable change in the booster mode, which is likewise reflected in the boost factor behavior. The results indicate: (i) the booster mode shifts to lower frequency, as seen from the displacement of the corresponding feature in the reflectivity and the shift of the boost factor line; (ii) the dip in reflected power broadens and deepens progressively as the tilt increases; (iii) the boost factor broadens and decreases correspondingly, (iv) the axion-induced field distribution is distorted and departs from the ideal TE_{11} shape



(a) Comparison of reflectivity from 3D FEM simulation and booster model.



(b) Comparison of boost factor from 3D FEM simulation and booster model.

Figure 5.9: Simulated (black dashed) and modeled (different colors solid) booster response under mirror tilt around the x -axis. The effective model parameters are fit to reproduce the reflectivities obtained from a 3D FEM simulation including mirror tilts from 0deg to 0.1deg. The spectra of reflected power $|\Gamma|^2$ and group delay τ_{gd} are shown in (a), whereas the boost factor spectra β^2 together with their relative difference are shown in (b). At the top of (b), the simulated transverse electric field distribution at the TE_{11} port is shown for the frequencies of maximum boost factor.

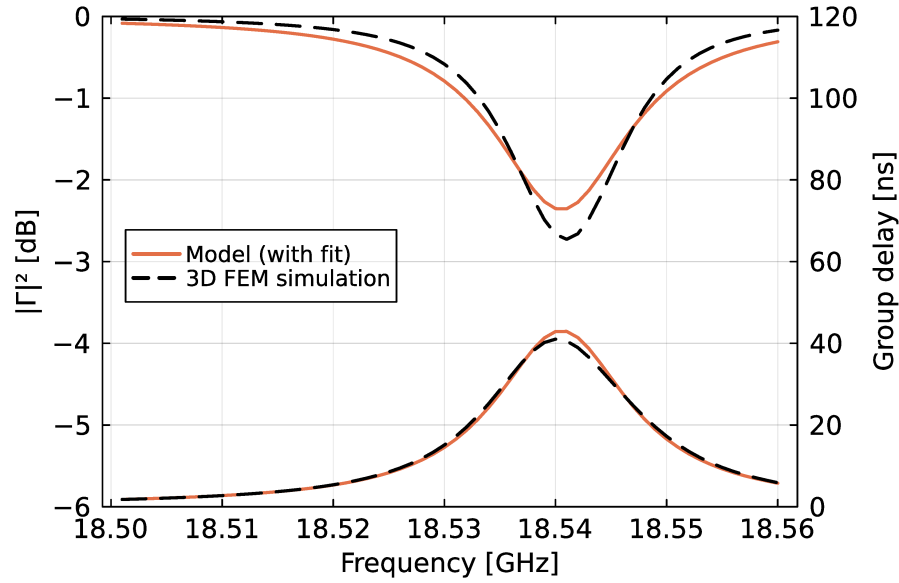
(case a).

The final step is to fit the booster model to the tilt simulation and derive the boost factor. The procedure described in section 5.2.1 is applied to reflectivity spectra from the 3D FEM simulation, that is disk positions d_i , disk phase thickness nd_{disk} and dielectric loss $\tan \delta$ of the booster model are fitted to match the simulated reflectivities. The boost factor that is computed from the fitted booster model is then compared to the simulation.

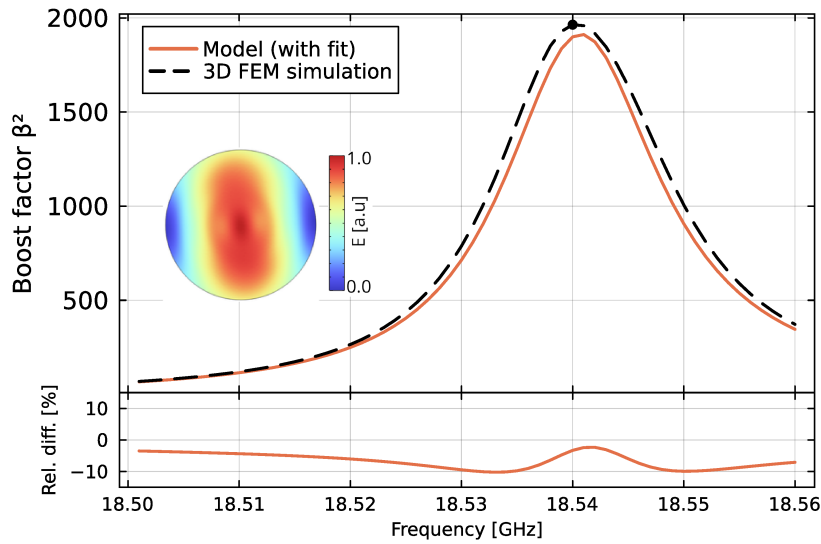
The results from the fitted model are overlayed with the simulated response in figure 5.9 (different colors, solid). The booster model reproduces the simulated reflectivities remarkably well in all cases except for the largest tilt of 0.1 deg. The modeled boost factor closely follows the frequency behavior of the simulation. Only the amplitude is increasingly overestimated for higher tilts. For a 0.05 deg tilt, which is compatible with the worst case of the runs later analyzed, the model still agrees with the 3D FEM simulation within 10 % at the boost factor maximum. At 0.1 deg, the boost factor maximum is overestimated by ~ 20 %. This behavior mainly arises from the degraded coupling to the axion current $\mathbf{J}_a$ not captured fully by the booster model: while the booster model uses its effective dielectric loss parameter to accurately reproduce the coupling visible in the reflected power $|\Gamma|^2$ as an increasingly deep dip, it underestimates the effect of the distorted electric field. This distortion, and hence decrease in overlap, limits the power that couples to the receiver and therefore reduces β^2 . This is a purely 3D effect beyond the scope of the booster model and demonstrates its limitations at large tilts. It was further checked that additionally rotating the tilted mirror around the z -axis

yields consistent results with the presented tilt orientation showing the highest deviation from the simulated boost factor.

Deviations of disk planarity (non-uniformity in the thickness across the surface of the disk) also modify the field distribution in a manner not captured by the booster model. Due to the localized resonance, we again focus on the region between mirror and closest disk. The measured planarity of that disk [85] is used to approximate a realistic disk shape within COMSOL with a min-max planarity of $\sim 50 \mu\text{m}$. It is included in the 3D FEM simulation of the ideal booster and boost factor and reflectivity simulations are repeated. The agreement of the model with the 3D FEM simulation after fitting its parameters is shown in figure 5.10. While the model does not fully reproduce the reflected power and group delay at the frequency of the booster mode, the modeled boost factor curve still closely follows the simulation and their maxima agree within 5 %. This also highlights the robustness of the booster model: although it does not reproduce the reflectivity exactly, it still manages to determine the boost factor within an acceptable uncertainty. Interestingly, the boost factor is underestimated in this case, showing that 3D effects can also lead to an increase in sensitivity.



(a) Comparison of reflectivity from 3D FEM simulation and booster model.



(b) Comparison of boost factor from 3D FEM simulation and booster model.

Figure 5.10: Comparison between reflectivity and boost factor obtained from the booster model (orange solid) and the 3D FEM simulation (black dashed) of a three-disk booster with a non-planar disk ($\sim 50\mu\text{m}$ min-max) in front of the mirror. The spectra of reflected power $|\Gamma|^2$ and group delay τ_{gd} are shown in (a), whereas the boost factor spectra β^2 together with their relative difference are shown in (b). The simulated transverse electric field distribution at the TE_{11} port is shown as an inset of (b) for the frequency of maximum boost factor.

In all cases, the largest deviation from the nominal value is seen in the effective dielectric loss parameter $\tan \delta$, which increased by about one order of magnitude. It is responsible for reproducing the decrease in boost factor amplitude. The shift in frequency is reproduced mainly by a deviation of nd_{disk} of the order of 10 % and d_3 of the order of 1 %. There is no qualitative difference between the effect of the mirror tilt and the effect of a deformed disk in the resulting effective parameters.

This analysis demonstrates that the booster model can account for the effect of small imperfections and 3D effects through the adjustment of effective disk positions, thicknesses, permittivities and losses. For the studied booster, the minimum value of $|\Gamma|^2$ at resonance is a reliable indicator for the deviation of the setup from the ideal case. Resonances with the minimum of $|\Gamma|^2$ above -5 dB are well reproduced by the model. This behavior is expected to be largely independent of the booster's transverse size, given a well-separated booster mode, which will be discussed below. The difference between simulated and modeled boost factor is by far dominated by a change in amplitude, visible by deviations of boost factor maximum of up to 10 %, as opposed to frequency. The study therefore confirms the uncertainty on $|\eta_A|^2$ of 12 % as estimated originally from electric field measurements [2].

5.3.3 Control of higher-order modes

The booster gains sensitivity with increasing radius. A key challenge is that this also allows propagation of an increased number of HOMs. These can be excited by imperfections and taper-induced mixing, then reflect within the volume and produce many additional resonance lines. An example of this overmoded behavior is illustrated in figure 5.11. Simulating this behavior accurately requires the use of the exact taper dimensions within the full-wave simulation. Performing a full 3D simulation becomes unfeasible at those scales and a 2D axis-symmetric simulation, as depicted in the top left of figure 5.12, is performed instead.

To justify the single-mode approximation of the booster model, the booster mode needs to be separated from HOMs in frequency. To illustrate our strategy, we include the results of a realistic simulation of the taper connected to the booster in figure 5.12. As can be seen from the included field maps (bottom left), the HOMs excited by the taper are mostly standing waves localized around the taper. Their frequencies are sensitive to the distance of the first disk to the taper aperture d_0 . In contrast, the electric field of the booster mode is mostly confined to the region containing the disks, determined by the disk spacings d_1 , d_2 , d_3 , and remains largely unaffected by d_0 . The modal map at the bottom-right clearly demonstrates this decoupling: the booster mode is constant at ~ 18.53 GHz, whereas the HOMs vary in frequency, depending on the distance of the leftmost disk to the taper aperture d_0 . The separation ring between taper and closest disk is therefore carefully

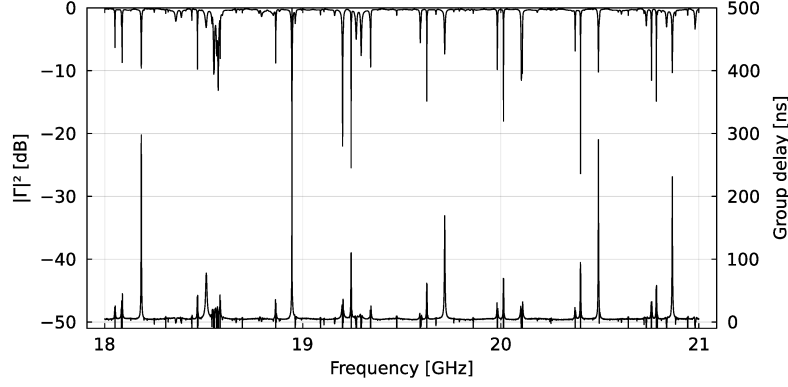


Figure 5.11: Exemplary reflectivity measurement showing reflected power $|\Gamma|^2$ (top curve) and group delay τ_{gd} (bottom curve) of CB200 over the full range of the taper, demonstrating the highly overmoded behavior of the booster.

chosen for each configuration to maximize the isolation between booster mode and HOMs in frequency.

Our tuning strategy was extensively tested and it is found to be very robust: in the case of CB200, a typical separation of HOMs of 50 MHz to 100 MHz is achieved across the taper’s full operational range of 18 GHz to 21 GHz. This separation justifies the single-mode approximation of the booster model and even allows further fine-tuning of the booster mode by externally offsetting the mirror. In this way, a scanning of ~ 50 MHz can be achieved without disassembling the booster and exchanging the separation rings.

5.3.4 Identification of the booster mode

Full-wave simulations of the booster are not practical for identifying the booster mode. As already mentioned, they fail to fully reproduce the measured reflectivity spectrum accurately because small imperfections and geometric details, such as those found in the dielectric lens, are particularly difficult to model. To unambiguously identify the booster mode, we measure the electric field distribution inside the booster volume. We employ a non-resonant, perturbation-based field mapping technique [86]. In this approach, $|\mathbf{E}|$ is locally inferred from changes in the measured reflectivity caused by a small dielectric object (bead) translated across the booster cross-section. The perturbed response is resolved by a VNA across the frequency range of interest, enabling direct correlation between the resonance lines in the reflectivity spectrum and their corresponding field profiles.

The setup is illustrated in figure 5.13, along with an exemplary field distribution measured for the CB200 configuration of run 1.1. The booster mode is distinguished by its lowest-order radial variation, characteristic for the TE_{11} dis-

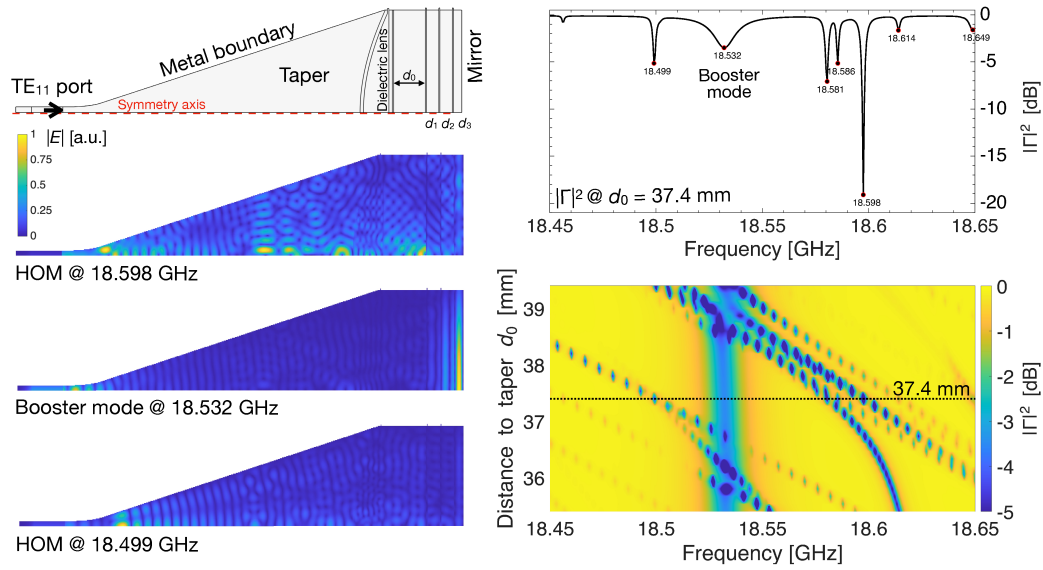


Figure 5.12: Axis-symmetric simulation of CB200 including the taper obtained for the booster configuration of run 1.1. (Top-left) Illustration of the geometry and boundary conditions of the simulated CB200 setup. The taper input allows only for TE_{11} propagation. (Bottom-left) Normalized E-field magnitude shown for the booster mode and adjacent HOMs at the indicated frequencies. (Top-right) Simulated reflectivity spectrum for $d_0 = 37.4$ mm, highlighting the booster mode and nearby HOMs. (Bottom-right) $|\Gamma|^2$ as obtained by a continuous sweep of the distance to taper d_0 .

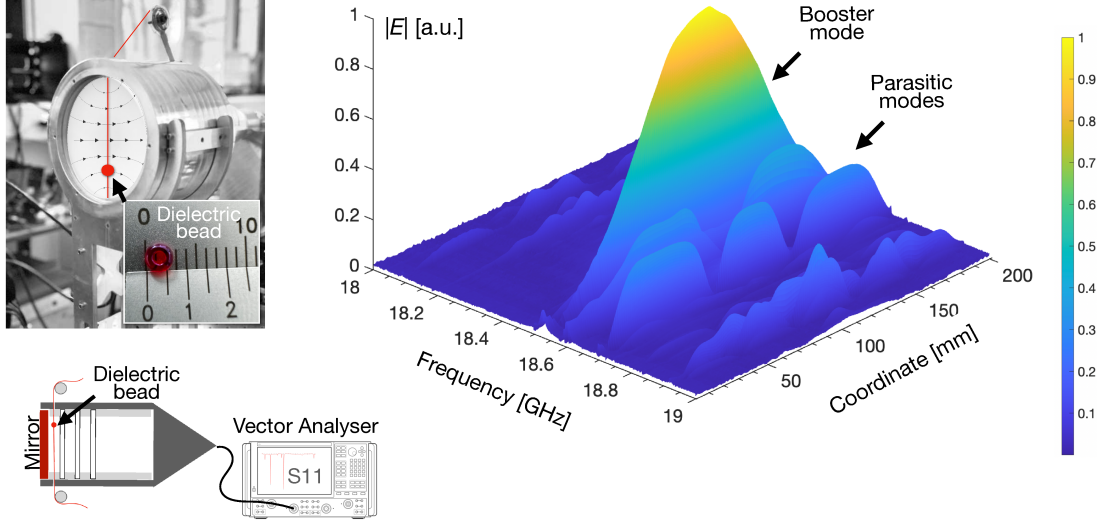


Figure 5.13: Perturbation-based field measurement implemented for the case of CB200, run 1.1. (Left) The E-field near the mirror surface is probed with a dielectric bead in the closed booster accessed by two sub-wavelength holes. The field is recorded as a function of the bead positions, indicated by the red line, using a VNA. The contour lines of the desired TE_{11} distribution are overlaid on the photo. (Right) $|E|$ measured along the line as indicated on the left of CB200 in the 18 GHz to 19 GHz range. The normalized field clearly reveals the radial structure of the modes. The booster mode appears as a well-defined parabolic shape consistent with the TE_{11} mode. Taken from supplementary material of [2].

tribution. In this qualitative approach, the field measurement is sufficient for identifying the booster mode without requiring absolute calibration, unlike [3, 4].

5.4 Sensitivity of a closed booster axion search

The CB200 MADMAX prototype, as shown in figure 6.1, was used to take data at CERN within the 1.6 T Morpurgo dipole magnet. Two different disk configurations were used, as described in the upper part of table 5.1. The configurations correspond to a maximum sensitivity to axion DM at around 18.55 GHz and 19.21 GHz respectively. By fine-tuning the mirror position, the boost factor maximum was shifted by $\mathcal{O}(10 \text{ MHz})$, enabling three physics-runs for configuration 1 and two in configuration 2 [2]. In this section it is described how the procedure is applied and how uncertainties are determined.

5.4.1 Receiver chain characterization

The power calibration of the receiver chain is performed as detailed in section 5.6.1 using a calibrated noise source with an excess-noise-ratio (ENR) of ~ 15 dB. The resulting calibration factors C (see equation (5.19)), converting uncalibrated power measurements to system temperature, are included in figure 5.18 in the section 5.6.1. For runs 1.2 and 1.3 the same power calibration was used and therefore an identical C is applied.

The uncertainty of C of around 10% is dominated by the uncertainty on the power output of the noise source and the attenuation factor of the attenuator connected between the noise source and the receiver. Further sources of uncertainty come from the physical temperature and the noise measurements themselves. The noise source includes a calibration certificate specifying frequency dependent ENR and associated uncertainty. The values are interpolated linearly between the given frequency points. For the attenuator, an attenuation of $A = 20.0(3)$ dB was approximated by the specified min-max difference of ~ 1 dB. The uncertainty on the physical temperature is conservatively assumed to be 5 K. The noise measurements are filtered by a third-order Savitzky-Golay filter [73] with a ~ 7 MHz window size. The uncertainty σ_N on the noise measurements is estimated by the mean squared error of the n bins of measurement N and filtered measurement N_{sg} :

$$\sigma_N = \sqrt{\frac{1}{n-1} \sum_i^n |N_{i,\text{sg}} - N_i|^2}. \quad (5.8)$$

All uncertainties are treated as uncorrelated and propagated into C .

The receiver chain noise is then modeled as described in section 5.2.2. The main source of parameter uncertainty stems from the power calibration. Since its uncertainty is almost fully correlated over frequency, a fit to the mean and $\pm 1\sigma$ band of the system temperature spectra is performed. This results in three sets of parameters that are then interpreted as mean and $\pm 1\sigma$ of the parameters themselves, defining a normal distribution for each model parameter.

The noise parameters determined by the fit are listed in table 5.3 in the section 5.6.2.

5.4.2 Boost factor determination

After verifying the frequency of the booster mode using the method demonstrated in section 5.3.4, the booster model is fit to the measured reflectivities of the different booster setups, as described in section 5.2.1. Figure 5.14 shows the agreement between modeled and measured reflected power $|\Gamma|^2$ and group delay τ_{gd} within the frequency range of the booster mode. The model reproduces all measurements

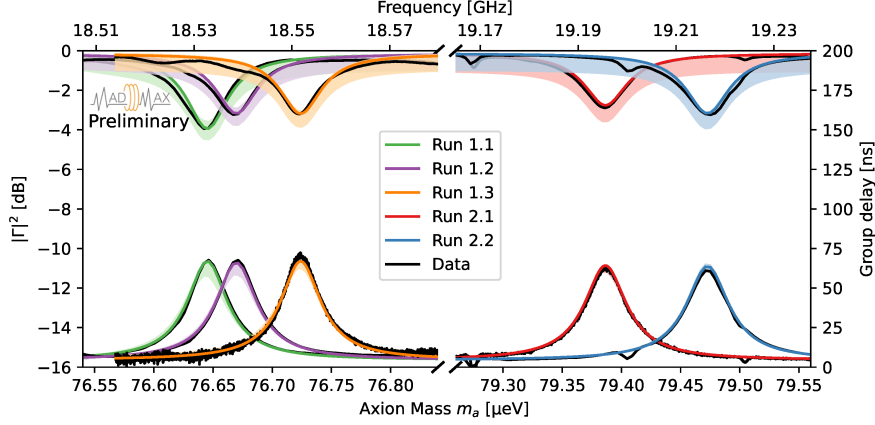


Figure 5.14: Reproduction of reflectivity measurements of the five runs of CB200 at the Morpurgo magnet (see table 5.1). The top curves show the reflected power $|\Gamma|^2$, while the bottom curves display the group delay τ_{gd} as a measure of phase. The band displays the uncertainty stemming from the fit procedure. The feature corresponding to the booster mode is reproduced within uncertainties for all runs.

within uncertainties, with the reflected power $|\Gamma|^2$ consistently being above -5 dB. This confirms the model’s assumption of a taper that is strongly coupled to the booster mode and is within the range for which the model was found to correctly predict the boost factor within the assigned uncertainty in section 5.3.2.

The resulting parameters and their fit uncertainties are listed in table 5.1. The biggest deviations between fitted and nominal value are seen for the effective dielectric loss parameter $\tan \delta$, which is one order of magnitude higher than the dielectric loss of sapphire, as well as for the disk phase thickness nd_{disk} , which is about 25 % higher than the nominal value. Since these parameters reflect the effective treatment of 3D effects, it is not expected that they match with the nominal values. For error propagation, a multivariate normal distribution is assumed, which is defined by the covariance matrix estimated from the fit result. The error band in figure 5.14 reflects the mean and $\pm 1\sigma$ of the fit uncertainty.

The system temperature spectra obtained with the combined noise model, using the booster parameters from table 5.1 and receiver noise parameters, as shown in table 5.3 in the section 5.6.2, are then fit to the measured system temperature spectra of the different booster setups. The results are shown in figure 5.15. Only the distance between LNA and booster L_{con} is left as a free fit parameter. The uncertainty band stems from propagating the uncertainties of the LNA and booster parameters using a Monte Carlo method, described further in section 5.4.3. The broad oscillation pattern, stemming from a standing wave between LNA and booster, as well as the resonance peak of the booster mode are reproduced for all

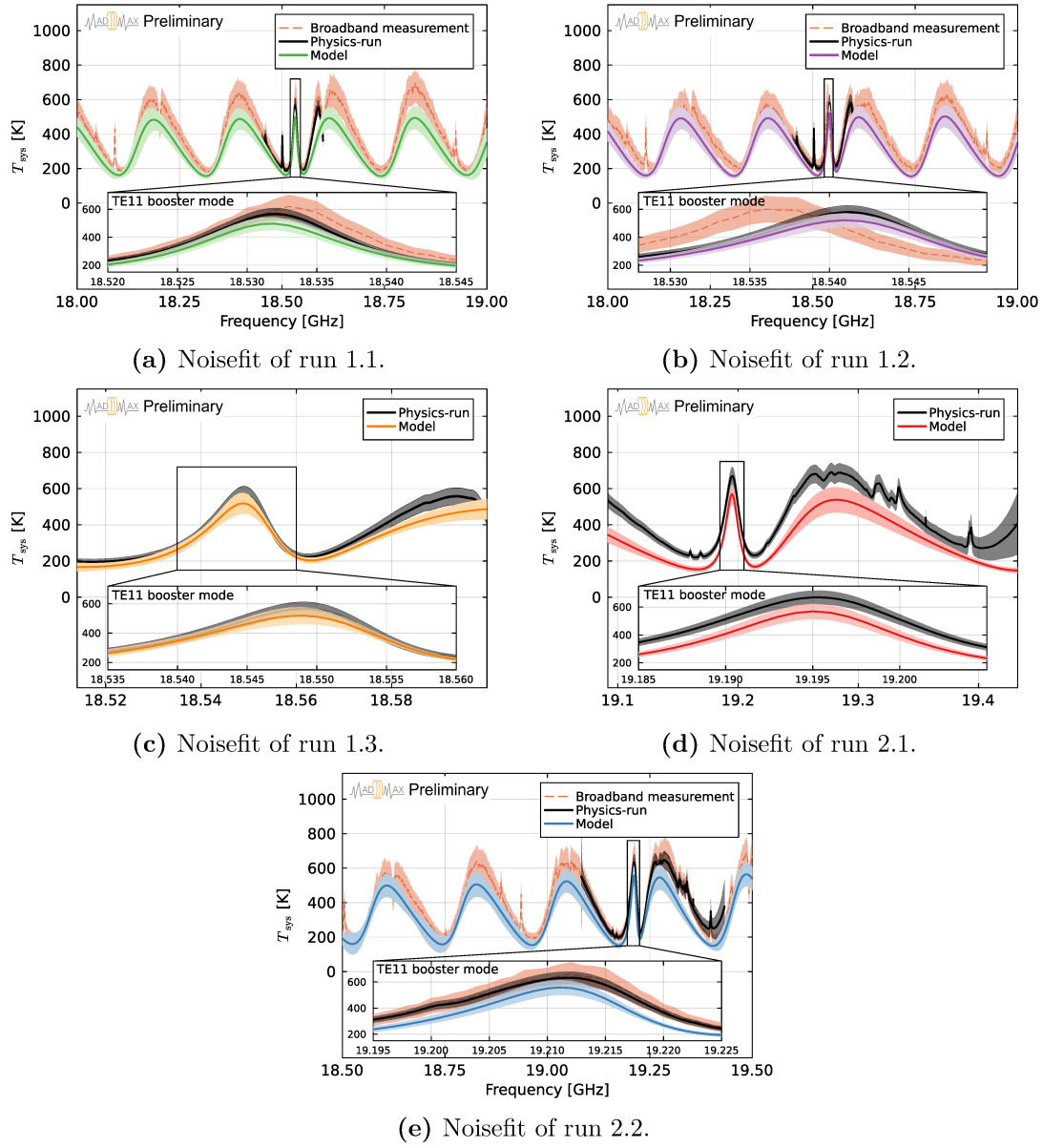


Figure 5.15: Reproduction of system temperature spectra of the five runs of CB200 at the Morpurgo magnet (see table 5.1). The noise model (different colors) matches the oscillation of the broadband measurement (orange) and the peak position of the booster mode of the physics-run (black), magnified in the inset. The bands signify the corresponding uncertainties. For runs 1.3 and 2.1, no broadband measurements are available. Figure 5.15a has already been shown as figure 5.7.

	d_1 [mm]	d_2 [mm]	d_3 [mm]	nd_{disk} [mm]	$\tan \delta$
Config 1	12.52(1)	12.25(1)	8.38(1)	3.06(4)	10^{-5}
Config 2	11.89(1)	12.25(1)	8.02(1)	3.06(4)	10^{-5}
Run 1.1	12.2(6)	12.14(8)	8.0979(2)	4.045(3)	$4.8(2) \times 10^{-4}$
Run 1.2	12.2(6)	12.14(7)	8.0953(2)	4.043(3)	$3.0(1) \times 10^{-4}$
Run 1.3	12.1(1)	12.13(2)	8.08954(4)	4.0402(7)	$3.00(3) \times 10^{-4}$
Run 2.1	11.72(5)	11.753(7)	7.80833(2)	3.8907(3)	$3.00(1) \times 10^{-4}$
Run 2.2	11.73(7)	11.752(9)	7.81706(2)	3.8981(4)	$2.00(1) \times 10^{-4}$

Table 5.1: Measured (top) and effective model (bottom) parameters. The effective parameters are extracted from a fit to the measured booster reflectivity. d_1 to d_3 describe the spacings between the disks, nd_{disk} is the product of refractive index n and disk thickness d_{disk} and $\tan \delta$ the (effective) dielectric loss of the sapphire disks.

setups within uncertainties, demonstrating the capability of the model to simulate the full system. For two of the runs, no broadband measurement exists.

In both cases, the previous run was performed using the exact same setup with only the mirror position changed using the tuning mechanism. Therefore, the best fit of the previous run was used as initial value for L_{con} in those cases. The fitted values are listed in table 5.4 in the section 5.6.2.

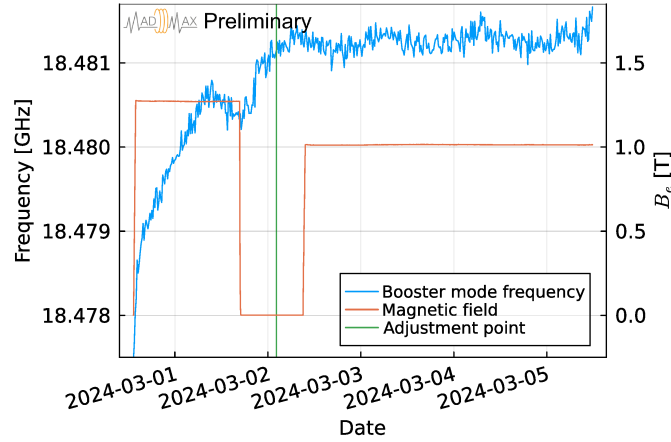


Figure 5.16: Booster mode frequency over time of run 1.2 (see table 5.1) magnet ramp up/down. The blue curve shows the significant increase of the resonance frequency over the first day. The orange curve displays the magnetic field B_e and the green bar signifies the time at which the first data point used in the analysis was taken.

In some cases, the booster mode resonance frequency had changed significantly over time, as evident in run 1.2. Figure 5.16 shows its evolution over the period

of data taking, with the booster mode frequency stabilizing only after a second ramp up of the magnet had finished. The main source of mechanical instability within the system is the mirror tuning mechanism, which is likely to take some time to stabilize after the initial setup. The only change apparent in the system temperature spectrum was a change of the resonance position of the booster mode, which mainly depends on the mirror position. In contrast, no change of frequency of higher-order-modes, that mainly depend on the distance of the first disk to the taper, is observed. It is therefore concluded that the change in resonance frequency is caused by a mechanical change of the mirror position. An additional single parameter fit of the mirror position within the combined booster and receiver model is performed to reproduce the frequency of the booster mode resonance in the system temperature spectrum, resulting in $\mathcal{O}(\mu\text{m})$ changes from the initially determined parameter. The modeled spectra shown in figure 5.15 already include this readjustment, which causes them to always match the resonance frequency of the physics-run, not necessarily the broadband calibration measurement. Data taken before the point of readjustment were discarded in the further analysis.

For the remaining time of the data taking the resonance frequency did not change significantly, with the standard deviation of the peak position being of $\mathcal{O}(100 \text{ kHz})$. The variation is determined for each run r individually and is assigned a normal distribution ΔF_r given by the mean and standard deviation of the frequency of the booster mode resonance, which is propagated to β^2 as uncertainty.

5.4.3 Uncertainty propagation

The total uncertainties on β^2 are dominated by the uncertainties of the booster model parameters and the field shape uncertainty. The field shape uncertainty of $|\eta_A|^2 = 0.84(10)$ is estimated from electric field measurements and confirmed by the simulations described in section 5.3. The uncertainties of the booster model parameters are given by a multivariate normal distribution estimated by the covariance matrix of their fit result. Further considered are the uncertainties of the receiver noise model parameters, given by individual normal distributions as estimated from the power calibration uncertainty. With all model parameters being assigned a distribution, a Monte Carlo method is employed to propagate the uncertainties to the final boost factor: The fits of L_{con} and d_3 to the measured system temperature are repeated $N = 100$ times, each time fixing all other parameters by drawing a random sample from their assigned distributions. Each resulting boost factor β_{ID}^2 is then further multiplied by a number drawn from the normal distribution 0.84 ± 0.1 to take into account the field shape uncertainty and shifted by a frequency drawn from ΔF_r to consider the observed frequency instability of the booster mode resonance. The result is considered a sample of the final boost factor distribution $\beta^2(\nu)$. This method results in N boost factor spectra, from which the

mean and standard deviation are determined for each frequency point. The mean boost factor varied by $< 1\%$ over the last 25 samples, confirming convergence.

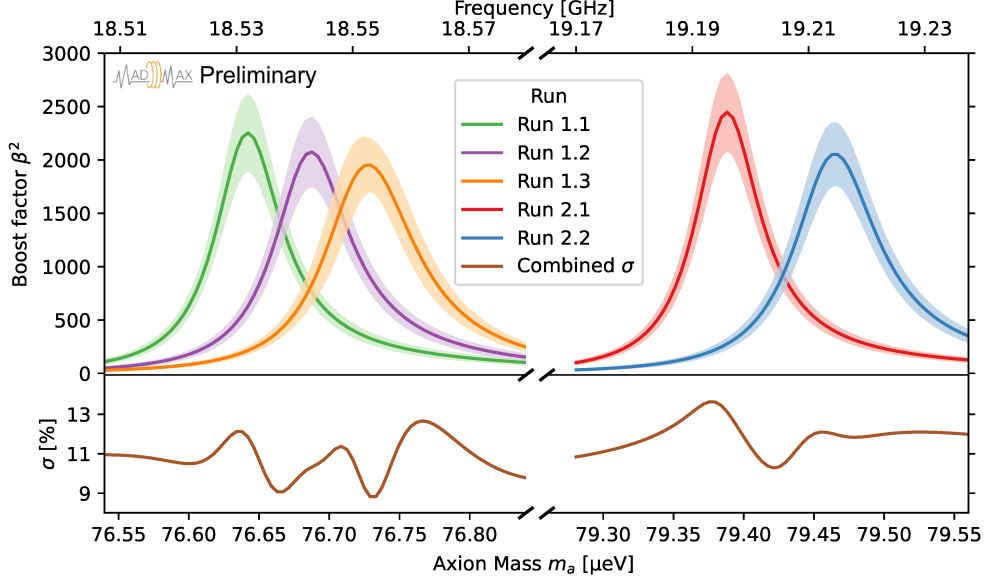


Figure 5.17: Boost factor distributions for the five runs (see table 5.1) and their combined relative uncertainty. The error band shows the standard deviation of the individual boost factor samples at each frequency point. The combined uncertainty shown in the bottom panel is the relative uncertainty of the boost factor sum with uncertainties added in quadrature.

For the determination of the sensitivity on $|g_{a\gamma}|$, the boost factors of all datasets are combined, assuming uncorrelated uncertainties. The combined relative uncertainty is shown in figure 5.17, together with the full boost factor distributions. The setup achieves boost factors of up to $\sim 2500(300)$ with a combined uncertainty between 9% to 14%, depending on frequency. With the boost factor distributions in place, a statistical method described in detail in [70], adapted from [87], was used to set limits on the axion parameter space in [2].

5.5 Conclusion

A method to determine the sensitivity of a closed dielectric haloscope to axion dark matter is described in this paper. A well-known transmission line approach is utilized to model the electromagnetic response of the setup. By fitting the model parameters, it is able to accurately reproduce the reflectivity of the booster, as measured by a VNA, and the system temperature spectra, as measured by

the receiver system. The model is validated against finite element simulations and is shown to absorb three-dimensional effects such as disk shape deviations and mirror tilt into effective parameters. It can therefore be used to determine the sensitivity to axion DM of the full setup, consisting of booster and receiver system. Uncertainties resulting from the power calibration, the fit procedure, the field shape and the time stability of the experiment are considered. The procedure is applied to data taken with the CB200 prototype at the Morpurgo magnet at CERN [2].

The presented model is able to replace time-consuming and computationally expensive simulations by simple calculations. The demonstrated ability to reproduce system temperature spectra makes it possible to relate system changes during data taking to specific model parameters. With this capability, unexpected changes of the setup can be adjusted for without interrupting data taking. Being largely volume independent, the model allows for further upscaling of the experiment. This work therefore lays the foundation for future dark matter searches by MADMAX, utilizing more complex booster systems with a higher number of disks and more extensive tuning capabilities.

Acknowledgments

The authors would like to thank the CERN central cryogenic laboratory and the CERN magnet team for the support during all the measurements with the Morpurgo magnet. We thank O. Reimann, D. Strom, and A. Hambarzumjan for significant contributions during the first phase of the MADMAX project. We thank O. Rossel for providing the drawing of the setup and artistic support. We acknowledge support by the Helmholtz Association (Germany); Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) under Project No. 550641633 and Germany's Excellence Strategy—EXC 2121 “Quantum Universe”—390833306. This work was also supported by the co-funded project HALOX between Agence national de la Recherche (ANR) and DFG grant (ANR-24-CE92-0049, DFG project No545526387). Computations were performed with computing resources granted by RWTH Aachen University under Project No. 7936. S. Knirck is supported by Fermi Research Alliance, LLC under Contract No. DEAC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics. We acknowledge the support of the MADMAX project by the Max Planck Society.

5.6 Appendix

5.6.1 Receiver power calibration

To take measurements of system temperature that are then reproduced by the noise model, the receiver system needs to be calibrated. The goal of this calibration is to convert power measurements performed with the receiver system to system temperature T_{sys} - that is the temperature at which a black body would emit a noise power equivalent to that of the connected device together with the receiver.

The components of the receiver system are shown at the bottom of figure 6.1. The booster is connected via a coaxial connection to a chain of amplifiers and filters in front of a spectrum analyzer (SA) to digitize the signal. The amplification is required, since the expected signal power is below the noise floor of a typical SA. The bandpass filters are added to avoid saturation of the amplifiers. A cavity filter at the end of the chain prevents image frequencies, mixer sidebands and aliasing effects from entering the data.

In practice, calibrating the receiver means determining its frequency dependent gain and added noise. This is done by connecting it to a calibrated noise diode, which is matched to the 50Ω coaxial connection between the receiver and the diode. Two measurements are performed, N_{on} and N_{off} , one with the noise source switched on and one with it switched off. It is assumed that the receiver chain behaves linearly within this range. The power emitted by the noise source P_N is given by the Rayleigh-Jeans approximation of Planck's law, that is

$$P_{N,\text{off}} = k_B B T_0 \quad (5.9)$$

$$P_{N,\text{on}} = k_B B 10^{\text{ENR}/10} T_0, \quad (5.10)$$

with the Boltzmann constant k_B , the resolution bandwidth of the receiver B , the physical temperature T_0 and the excess noise ratio (ENR) of the noise source. To not saturate the receiver chain, an attenuator of $A = 20\text{ dB}$ is added between receiver chain and noise source, resulting in a modification to equation (5.10):

$$P_{N,\text{on}} = k_B B [10^{(\text{ENR}-A)/10} T_0 + (1 - 10^{-A/10}) T_0] \equiv k_B B T_N, \quad (5.11)$$

where the term in square brackets is the sum of the attenuated noise source's power and the thermal noise emitted by the attenuator itself.¹ The receiver chain also adds noise to the measurements, quantified by its equivalent noise temperature T_e . It is assumed to be fully uncorrelated to the noise emitted by the noise source, which would result in a simple sum for the total power $P_{\text{on/off}} = k_B B T_e + P_{N,\text{on/off}}$.

¹From this, equation (5.9) can be recovered by setting $\text{ENR} = 0$ and equation (5.10) by setting $A = 0$

In the present case, however, the receiver chain is not perfectly matched to the 50 Ω coaxial connection, resulting in some of the noise source's power being reflected. Therefore, the actual total power delivered to the receiver chain is [80]

$$P_{\text{on/off}} = k_B B T_e + (1 - |\Gamma_{\text{RC}}|^2) k_B B T_{\text{N}/0}, \quad (5.12)$$

with the receiver's input port reflection given by Γ_{RC} . The factor can be understood intuitively by thinking about the reflected power corresponding to $|\Gamma_{\text{RC}}|^2$ being subtracted from the total power.

Note that this simple expression is only possible due to the fact that the noise diode itself is matched to the coaxial connection, which is not true for the booster.

For the matched noise diode, the receiver chain gain G is given by the measured noise

$$N_{\text{on/off}} = G P_{\text{on/off}}, \quad (5.13)$$

which, by inserting equation (5.12), can be written as

$$N_{\text{on}} = G k_B B [T_e + (1 - |\Gamma_{\text{RC}}|^2) T_{\text{N}}] \quad (5.14)$$

$$N_{\text{off}} = G k_B B [T_e + (1 - |\Gamma_{\text{RC}}|^2) T_0]. \quad (5.15)$$

The term in brackets defines the system temperature T_{sys} for a given device-under-test (DUT) with equivalent temperature T_{DUT} :

$$T_{\text{sys}} \equiv T_e + (1 - |\Gamma_{\text{RC}}|^2) T_{\text{DUT}}. \quad (5.16)$$

Solving equation (5.14) and equation (5.15) for G and T_e results in

$$G = \frac{N_{\text{on}} - N_{\text{off}}}{P_{\text{on}} - P_{\text{off}}} \quad (5.17)$$

$$T_e = \frac{N_{\text{on/off}}}{G k_B B} - (1 - |\Gamma_{\text{RC}}|^2) T_{\text{N}/0}. \quad (5.18)$$

The resulting G not only includes the gains and attenuations of the different filters and amplifiers of the receiver system but also the insertion loss of the mixer and eventual internal gain from the analog to digital converter used to digitize the signal. The equivalent noise temperature should roughly match the specified datasheet value of the first stage amplifier, since all other noise sources are suppressed by the first stage amplifier's gain. To then convert any measured power spectrum N to system temperature T_{sys} , it only needs to be multiplied by the calibration factor

$$C = \frac{1}{G k_B B} = \frac{P_{\text{on}} - P_{\text{off}}}{k_B B (N_{\text{on}} - N_{\text{off}})} = (1 - |\Gamma_{\text{RC}}|^2) \frac{T_{\text{N}} - T_0}{N_{\text{on}} - N_{\text{off}}}. \quad (5.19)$$

The resulting calibration factors C for the five runs (see table 5.1) together with their uncertainties are shown in figure 5.18.

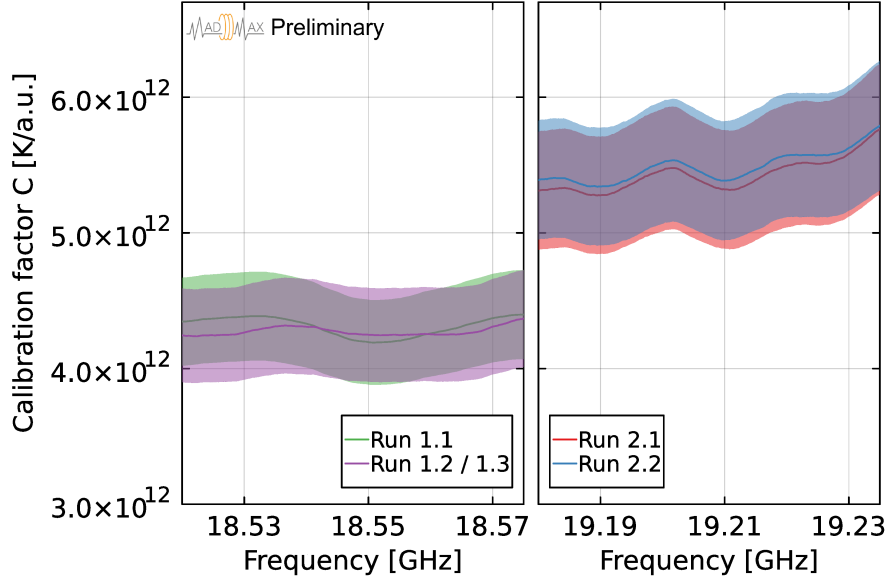


Figure 5.18: Calibration factors of the five runs of CB200 at the Morpurgo magnet (see table 5.1) uncertainties. Three runs were performed around 18.55 GHz (left) and two around 19.21 GHz (right). Note that runs 1.2 and 1.3 use the same power calibration.

5.6.2 Model parameters

Besides the fitted parameters, the booster model requires a set of fixed parameters listed in table 5.2. $\epsilon_d = n^2$ and d_{disk} are included in the fit in the form of their product nd_{disk} . d_0 and σ_{Al} are kept fixed throughout. Since d_0 does not affect the boost factor calculation, its uncertainty does not need to be considered further. While no clear uncertainty on σ_{Al} was found in the literature, it was checked that a 20% variation only changes the boost factor by $< 2\%$, therefore its uncertainty is neglected.

The fit parameters of the receiver noise model and combined model are shown in table 5.3 and table 5.4 respectively.

$L + d_0$ [mm] (C1)	$L + d_0$ [mm] (C2)	d_{disk} [mm]	ϵ_d	σ_{Al} [S m^{-1}]
203(1)	187(1)	1.00(1)	9.36(10)	3.77×10^7

Table 5.2: Fixed parameters used in the booster model for configuration 1 (C1) and configuration 2 (C2).

V_n [pV]	I_n [pA]	$ c $	ϕ_c	L_{open} [ps]	L_{short} [ps]
435(30)	17(1)	0.6(1)	358.7(9)	45.8(1)	71.54(8)

Table 5.3: Fitted receiver noise model parameters, consisting of voltage and current noise amplitude V_n and I_n , their correlation magnitude and angle c and ϕ_c as well as the electrical lengths L_{open} and L_{short} between open and short standard and LNA respectively.

	L_{con} [ns]	d_3 [mm]
Run 1.1	1.58(7)	8.097(1)
Run 1.2	1.55(5)	8.0916(5)
Run 1.3	1.78(47)	8.0882(6)
Run 2.1	1.50(4)	7.808 33(2)
Run 2.2	1.22(4)	7.8155(1)

Table 5.4: Fit values for the distance between LNA and booster in the noise model L_{con} and the readjustment of the distance between mirror and closest disk d_3 .

5.7 Additional material

5.7.1 Analytical expression of the noise model

The presented study uses the commercial software ADS to implement the noise model of the receiver system. This section derives an analytical equation, by treating the voltage and current noise sources as coherent signals. The goal is to simulate the power delivered to the LNA, given by

$$P = \frac{1}{2} \operatorname{Re}\{V_{\text{LNA}} \overline{I_{\text{LNA}}}\} = \frac{1}{2} \frac{|V_{\text{LNA}}|^2}{\operatorname{Re}\{Z_{\text{LNA}}\}}. \quad (5.20)$$

Voltages and currents are split into right and left moving amplitudes, as per Fig. 5.4,

$$V = V^+ e^{-i\beta z} + V^- e^{i\beta z} \quad (5.21)$$

$$I = I^+ e^{-i\beta z} + I^- e^{i\beta z}, \quad (5.22)$$

that are propagating with propagation constant β along the z-axis. Furthermore,

$$V^+ = \frac{I^+}{Z_L} \quad (5.23)$$

$$V^- = -\frac{I^-}{Z_L} \quad (5.24)$$

at a load of impedance Z_L . A voltage noise source of amplitude V_n corresponds to right and left moving voltages with $V_{in}^+ = 1/2 V_n$ and $V_{in}^- = -1/2 V_n$. The current noise source of amplitude I_n is connected in parallel, resulting in $I_{in}^+ = 1/2 I_n$ and $I_{in}^- = 1/2 I_n$. To consider correlations between voltage and current source, the current is split into a correlated and uncorrelated part by $I_n = c I_n + (1 - c) I_n = I_{n,c} + I_{n,u}$ with $c \in \mathbb{C}$. While the correlated part gets added to the total voltage coherently, the uncorrelated part is later added incoherently to the measured power. Now consider any DUT with reflectivity Γ_{DUT} directly connected to the LNA. Propagating waves in direction of the DUT are reflected towards the LNA and waves propagating towards the LNA are reflected towards the DUT, resulting in a sum of reflected fields given by the geometric series

$$E = E_{in} \cdot (1 + \Gamma_{\text{DUT}} \Gamma_{\text{LNA}} + \Gamma_{\text{DUT}} \Gamma_{\text{LNA}} \cdot \Gamma_{\text{DUT}} \Gamma_{\text{LNA}} + \dots) \quad (5.25)$$

$$= E_{in} \cdot \left(\sum_{k=0}^{k=\infty} (\Gamma_{\text{DUT}} \Gamma_{\text{LNA}})^k \right) \quad (5.26)$$

$$= \frac{E_{in}}{1 - \Gamma_{\text{DUT}} \Gamma_{\text{LNA}}} \equiv E_{in} \cdot \Gamma_{tot}. \quad (5.27)$$

Any eventual transmission line between LNA and DUT can easily be absorbed into Γ_{DUT} . The right and left moving parts of voltages and currents can now be written as:

$$V^+ = V_{in}^+ \cdot (\Gamma_{tot} - 1) + V_{in}^- \cdot \Gamma_{\text{LNA}} \Gamma_{tot} \quad (5.28)$$

$$V^- = V_{in}^- \cdot \Gamma_{tot} + V_{in}^+ \cdot \Gamma_{\text{DUT}} \Gamma_{tot} \quad (5.29)$$

$$I^+ = I_{in}^+ \cdot (\Gamma_{tot} - 1) + I_{in}^- \cdot \Gamma_{\text{LNA}} \Gamma_{tot} \quad (5.30)$$

$$I^- = I_{in}^- \cdot \Gamma_{tot} + I_{in}^+ \cdot \Gamma_{\text{DUT}} \Gamma_{tot}. \quad (5.31)$$

Since in case of a perfectly matched DUT the power delivered to the LNA would only correspond to the left moving waves, the first term of the geometric series is subtracted from V_{in}^+ and I_{in}^+ in the sum of right propagating fields. Because the fields enter between LNA and DUT, some terms are first multiplied by Γ_{DUT} or Γ_{LNA} . The total voltage at the LNA due to correlated voltage and current noise is now given by

$$V_{\text{LNA},c} = V^+ + V^- + Z_{\text{LNA}}(I_{n,c}^+ - I_{n,c}^-) \quad (5.32)$$

$$= \frac{1}{2} (V_n(\Gamma_{tot} - 1) - Z_{\text{LNA}} c I_n(\Gamma_{tot} + 1)) \quad (5.33)$$

and the voltage due to uncorrelated current noise by

$$V_{\text{LNA},u} = Z_{\text{LNA}}(I_{n,u}^+ - I_{n,u}^-) \quad (5.34)$$

$$= -\frac{1}{2} Z_{\text{LNA}}(1 - c) I_n(\Gamma_{tot} + 1). \quad (5.35)$$

The total power is then

$$P_{\text{LNA}} = \frac{1}{2} \left(\frac{|V_{\text{LNA},c}|^2 + |V_{\text{LNA},u}|^2}{\text{Re}\{Z_{\text{LNA}}\}} \right). \quad (5.36)$$

In some cases the DUT might as well emit some noise towards the LNA. This noise is considered fully uncorrelated to the LNA noise and therefore treated separately. Analogous to Eq. 5.28 to Eq. 5.31, one can write the following, considering now the DUT to be a generator itself with $V_{in}^- = 1/2 V_{\text{DUT}} e^{-i\omega d}$ and $V_{in}^+ = 0$ (since the generator is matched to the DUT) with d the distance between DUT and LNA in units of time:

$$V^+ = V_{in}^- \Gamma_{\text{LNA}} \Gamma_{tot} \quad (5.37)$$

$$V^- = V_{in}^- \Gamma_{tot}. \quad (5.38)$$

Just like before, Γ_{DUT} is assumed to include the phase term of the transmission line between the devices. From this, it follows that the power delivered to the

LNA due to the DUT emissions is

$$P_{\text{DUT}} = \frac{1}{8} \operatorname{Re} \left\{ V_{\text{DUT}} \Gamma_{\text{tot}} (1 + \Gamma_{\text{LNA}}) \cdot \overline{V_{\text{DUT}} \Gamma_{\text{tot}} \frac{1 + \Gamma_{\text{LNA}}}{Z_{\text{LNA}}}} \right\} \quad (5.39)$$

$$= \frac{1}{8Z_0} |V_{\text{DUT}}|^2 \frac{1 - |\Gamma_{\text{LNA}}|^2}{|1 - \Gamma_{\text{DUT}} \Gamma_{\text{LNA}}|^2} \equiv \frac{1}{8Z_0} |V_{\text{DUT}}|^2 F_{\text{LNA}}, \quad (5.40)$$

which differs from the emitted power by the factor F_{LNA} . This equation holds true in general for any sort of power emitted by the DUT. A typical scenario is that of a passive device which emits thermal radiation corresponding to its emissivity. In case of a perfect black body with resistance R and temperature T , the emitted thermal radiation is given by the microwave approximation of Planck's law:

$$V_n = \sqrt{8k_B T R}, \quad (5.41)$$

assuming a noise bandwidth of 1 Hz for simplicity. For $R = Z_0$, as expected this yields a power delivered to the LNA of $P_{\text{DUT}} = k_B T$. More generally, the emissivity of a passive device can be approximated by its absorption, which, in case of no dissipated power, is given by the reflectivity as $1 - \Gamma_{\text{DUT}}^2$. By combining this with Eq. 5.36 and Eq. 5.40 and inserting Eq. 5.33 and Eq. 5.35, the full power spectrum of any passive DUT connected to a noisy LNA is given by:

$$P = \frac{1}{8 \operatorname{Re}\{Z_{\text{LNA}}\}} (|V_n(\Gamma_{\text{tot}} - 1) - Z_{\text{LNA}} c I_n(\Gamma_{\text{tot}} + 1)|^2 + |Z_{\text{LNA}}(1 - c) I_n(\Gamma_{\text{tot}} + 1)|^2) + (1 - |\Gamma_{\text{DUT}}|^2) F_{\text{LNA}} k_B T \quad (5.42)$$

6 Search for axion dark matter using a closed booster

This section is published as *"First Search for Axion Dark Matter with a MAD-MAX Prototype"* [2]. In section 6.6, additional unpublished material is presented.

Significance: This publication demonstrates the practical application of the previously developed procedure to determine the boost factor of a closed dielectric haloscope. It proves the viability of the dielectric haloscope concept for axion searches at masses around $100 \mu\text{eV}$, including limited tuning capabilities. The determination of the boost factors presents the main contribution by the thesis author.

6.1 Introduction

Axions have become prominent candidates for cold dark matter (DM) [88, 89, 90, 91]. They were originally introduced to explain the absence of CP-violating effects in quantum chromodynamics (QCD) through the Peccei-Quinn (PQ) mechanism [75, 11, 76, 12]. In the well motivated post-inflationary PQ-symmetry breaking

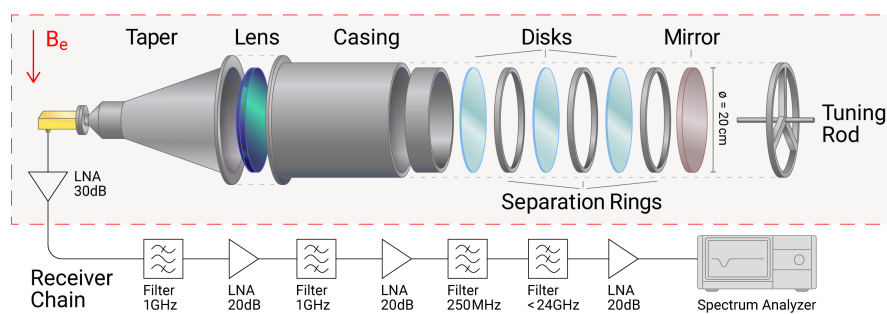


Figure 6.1: Exploded schematic view of the MADMAX prototype CB200 and the receiver chain. The shaded region represents the components exposed to the B-field.

scenario the axion mass required to match the observed DM density is expected

to be in the range $m_a \sim 26$ to $500 \mu\text{eV}$ for axion models with short-lived domain walls [45, 92, 15, 16] or even higher for other models [93, 94]. In the following, the term axion will refer to both, QCD axions and axion-like particles.

In an external magnetic field $\mathbf{B}_e$, axions convert to photons and source an oscillating current with frequency corresponding to m_a , given by $\mathbf{J}_a = g_{a\gamma}\mathbf{B}_e\dot{a}$ [48] as an additional term to be added to Maxwells' equations, with $g_{a\gamma}$ the axion-photon coupling strength and $\dot{a}$ the temporal derivative of the axion field. Cavity searches are a promising approach to detect axions [74]. They are designed to convert axions from the DM halo into microwave photons using a resonator in a magnetic field. Several cavity searches [58, 95, 96, 97, 59, 98, 99, 87, 60, 100] have already put significant constraints on $|g_{a\gamma}|$ for $1 \mu\text{eV} \lesssim m_a \lesssim 26 \mu\text{eV}$, with first efforts being underway to test higher m_a [61, 101, 102, 103, 104, 105, 65, 106, 107, 108], which are challenging to probe with cavity experiments as their size must be adjusted to match the wavelength of the generated photon. At higher m_a values, corresponding to higher frequencies of the generated photon, this results in a reduced volume for resonant axion-photon conversion, thereby decreasing sensitivity.

The MAgnetized Disk and Mirror Axion eXperiment (MADMAX) [79, 23, 24] makes use of the dielectric haloscope concept to access the well-motivated mass range $m_a > 26 \mu\text{eV}$. This concept, based on the magnetized mirror axion DM search idea [78], uses a *booster* composed of a set of parallel dielectric disks and a metallic mirror to enhance a potential axion signal, effectively decoupling the conversion volume from the wavelength and thereby overcoming the limitation of typical cavity experiments towards higher masses.

Due to the large de Broglie wavelength of the considered non-relativistic DM axions of $\mathcal{O}(10 \text{ m})$, $\mathbf{J}_a$ is assumed to be uniform over the extent of the booster. This leads to an oscillating electric field uniform across a single medium with dielectric permittivity ϵ . The field $\mathbf{E}_a$ exhibits discontinuities at the boundaries of different media, resulting in the emission of traveling waves with frequency corresponding to m_a ensuring the continuity of the electric field. The signal power P_{sig} due to the coherent emission from multiple surfaces in relation to the power emitted by an ideal magnetized mirror, P_0 , is enhanced by the frequency-dependent boost factor $\beta^2 = P_{\text{sig}}/P_0$. By adjusting the disk positions its center frequency and width can be tuned. This makes it possible to boost the potential DM axion signal compared to the magnetized mirror concept [78] in dedicated mass ranges. The sensitivity of a MADMAX prototype setup to $|g_{a\gamma}|$ for a specific signal to noise ratio (SNR)

can be estimated as [48]

$$\begin{aligned}
 |g_{a\gamma}| &= 3.5 \times 10^{-11} \text{ GeV}^{-1} \sqrt{\frac{2 \times 10^3}{\beta^2}} \sqrt{\frac{T_{\text{sys}}}{300 \text{ K}}} \\
 &\times \left(\frac{0.1 \text{ m}}{r} \right) \left(\frac{1 \text{ T}}{B_e} \right) \left(\frac{2.2 \text{ days}}{\Delta t} \right)^{1/4} \sqrt{\frac{\text{SNR}}{5}} \\
 &\times \left(\frac{m_a}{80 \text{ } \mu\text{eV}} \right)^{5/4} \sqrt{\frac{0.3 \text{ GeV/cm}^3}{\rho_a}},
 \end{aligned} \tag{6.1}$$

with system temperature T_{sys} , radius r of the disks and mirror, effective data-taking time Δt , and the local axion DM density ρ_a .

The MADMAX Collaboration has made significant progress in advancing the development of a future large-scale dielectric haloscope [109, 25, 67, 64, 4, 110, 111, 85]. Small-scale setups have successfully validated the dielectric haloscope approach in dark photon DM searches [112, 3], specifically with a broadband MADMAX prototype [3]. This work for the first time demonstrates that a tunable dielectric haloscope can be used to search for axion DM.

6.2 Experimental Setup

The Closed Booster 200 (CB200) MADMAX prototype is used. As illustrated in Fig. 6.1, it consists of an aluminum mirror and three sapphire disks of 1 mm thickness, each with a diameter of 200 mm and distanced by separation rings enclosed in an aluminum casing, a Rexolite[®] microwave lens and an aluminum taper, all controlled at $\mathcal{O}(10 \text{ } \mu\text{m})$ precision. In contrast to the design used in [3] and envisioned for the final MADMAX design, the prototype is enclosed by conductive boundaries, resulting in fixed boundary conditions, such that only a limited number of cylindrical wave-guide modes within a given frequency range exists inside the booster. This significantly simplifies modeling of the electromagnetic response and allows us to determine β^2 with fewer dedicated measurements. Approximately 84% of the axion induced power couples to the fundamental transverse electric TE_{11} mode as calculated by the modal overlap formalism considering the mode of propagation in a wave guide, which is absorbed in the calculation of β^2 . The emitted TE_{11} power emissions are coupled into the receiver chain by a lens and a specifically designed taper via a $50 \text{ } \Omega$ transmission line. The power output of the system is coupled to a heterodyne receiver system, consisting of a series of low noise amplifiers (LNAs) and filters, and a Rohde & Schwarz FSW43 real-time spectrum analyser (SA), which streams time-domain data with $\sim 0.11 \text{ s}$ coherent integration time to a computer where it is Fourier transformed on GPUs to power spectra.

Subsequently, 8047 single power spectra are averaged to individual files with corresponding $8047 \times 0.11 \text{ s} = 14.75 \text{ min}$ measurement time each. The receiver system has a bandwidth of $\sim 250 \text{ MHz}$, which is approximately centred around the respective β^2 peak, a resolution of 9 Hz and provides negligible dead-time. It was configured for each run to acquire data around the frequency of the expected maximum β^2 . The gain of the receiver chain was calibrated by the Y-factor method [80] using a calibrated noise diode. T_{sys} is dominated by the $\sim 230 \text{ K}$ added noise of the LNA. Additional information on the setup and the methodology used can be found in [113].

Two different disk spacing *configurations* leading to β^2 peaks in a $\sim 10 \text{ MHz}$ range were used to search for axions: Configuration 1 is designed to be sensitive around 18.55 GHz and configuration 2 around 19.21 GHz , corresponding to axion masses around $76.72 \mu\text{eV}$ and $79.45 \mu\text{eV}$, respectively. Frequency adjustment is done by proper choice of the width of separation rings between adjacent disks and between mirror-facing disk and mirror: 12.52 mm , 12.25 mm and 8.38 mm (from left to right in Fig. 6.1) for configuration 1 and 11.89 mm , 12.25 mm and 8.02 mm for configuration 2. The widths are optimized for a resonance of the TE_{11} mode between mirror and adjacent disk. A $\mathcal{O}(\mu\text{m})$ change in separation between mirror and closest disk results in $\mathcal{O}(\text{MHz})$ frequency shift of the β^2 distribution. The peak frequencies for both configurations were precisely tuned for different runs to different values by applying mechanical force to the mirror using an adjustable tuning rod. Three *physics-runs* around 18.55 GHz and two around 19.21 GHz were performed, sensitive to different frequency ranges with peak frequencies separated by $\sim 10 \text{ MHz}$.

Measurements of CB200 power spectra with a total measurement time of 14.5 days have been performed inside the Morpurgo dipole magnet [114] at CERN. The magnet was operated at field strengths from $1.00(1) \text{ T}$ to $1.58(2) \text{ T}$.

6.3 Boost factor determination

The boost factor β^2 , which is strongly correlated to the reflectivity S_{11} of the system [48], and its corresponding uncertainties are determined using a single-mode booster model [113] implemented in ADS [115]. The model is fitted to S_{11} measurements of the booster. Disk positions, thicknesses, dielectric losses and permittivity are extracted. Losses due to the lens and three-dimensional effects, such as tilts, are modeled by effective dielectric losses.

Identification of the TE_{11} booster mode is done for each configuration by measuring the shape of the E-field between mirror and closest disk, using a bead-pull method [4]. It follows the expected shape. S_{11} measurements confirm the separation of the TE_{11} modes from parasitic modes.

The β^2 calculation includes the overlap between the theoretical field shape and the uniform axion current of 84%. This is analogous to a form factor in cavity experiments. The uncertainty on the overlap of $\pm 12\%$ is determined from the deviation between expected and measured field shape obtained from bead-pull measurements.

When measuring power spectra, the experimentally determined $\sim 41.5 \Omega$ impedance of the first LNA of the receiver chain leads to a mismatch with the transmission line inducing a standing wave between booster and LNA. This modifies β^2 obtained from S_{11} measurements, depending on the distance between the booster and this LNA. The distance is determined from the broadband oscillation pattern induced by resonance of the LNA noise, similar as in [3]. The modeled frequency behavior of the power spectra, adapted to consider the LNA impedance mismatch, matches the measured ones from the physics-runs. Broadband power spectra measurements were taken prior to each physics-run. Drifts in T_{sys} or frequency are monitored by comparing the individual 14.75 min long measurements. The variation of the measured booster peak frequencies during physics-runs is used to evaluate the systematic uncertainty on β^2 . They follow a Gaussian distribution with a width around 200 kHz for all physics-runs, resulting in an additional uncertainty in β^2 for a given frequency of less than 4%.

The extracted boost factors including their uncertainties are shown in Fig. 6.2. The maximum values for β^2 are around 2000, matching expectations [48] and have uncertainties of 13 to 17%.

6.4 Analysis

The analysis procedure for the MADMAX axion DM search [70] is modeled after the framework established by HAYSTAC [116] and adopts their nomenclature. It combines a set of individual power spectra $P_i(\nu)$, each measured with $\mathcal{O}(14.75 \text{ min})$ averaging time, to a *grand spectrum*, where the full axion signal information is contained in a single bin. The raw FFT data are re-binned for analysis to 0.9 kHz, enough to resolve a potential ALP DM signal.

The spectra $P_i(\nu)$ are filtered using a fourth-order Savitzky-Golay (SG) filter [73] with a window length of $\sim 1 \text{ MHz}$ to extract the baselines $P_{\text{bl},i}(\nu)$, which are subsequently used to obtain the processed spectra $p_{\text{proc},i} = (P_i/P_{\text{bl},i} - 1)$. These processed spectra give local power excesses above the baseline across all frequency bins i . In the absence of an axion signal and assuming ideal baseline removal, the $p_{\text{proc},i}$ are expected to follow Gaussian white noise with zero mean and a standard deviation determined by the averaging time. Studies of the full analysis chain applied to simulated data including synthetic axion signals reveal an SNR attenuation due to the SG filter by a factor of $\eta_{\text{SG}} = 0.95$, which is constant in the range

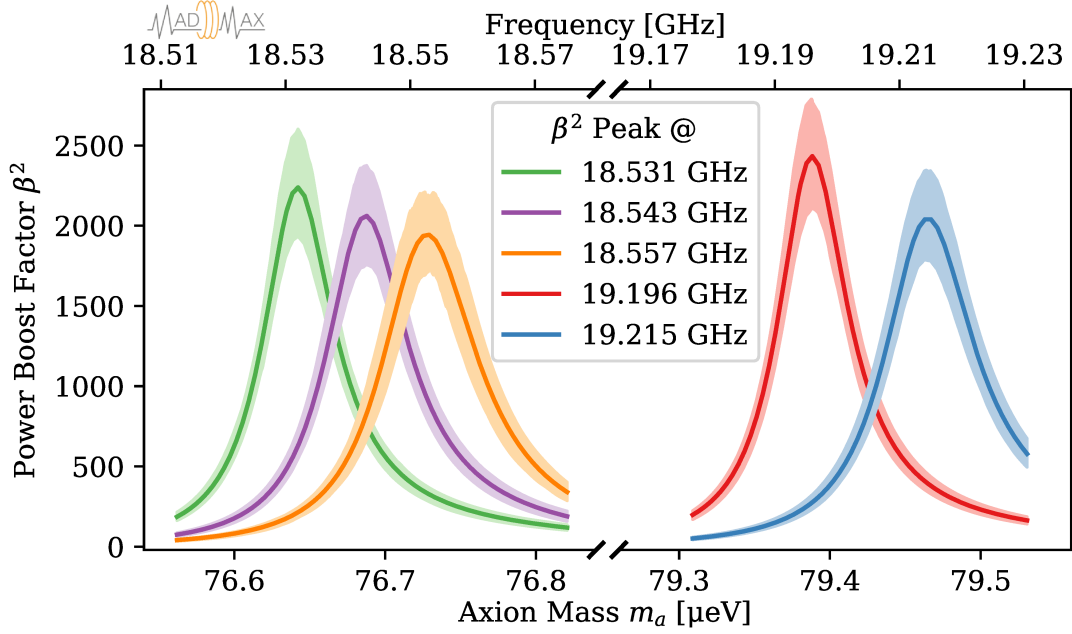


Figure 6.2: Modeled β^2 distributions for the five physics-runs. Lines denote mean β^2 and shaded regions give $\pm 1\sigma$ intervals.

$1 < \text{SNR} < 10$.

The baseline of the recorded power spectra unexpectedly fluctuates on the scale of ~ 9 kHz which introduces unwanted correlations between the processed spectra. This sinusoidal fluctuation is removed by applying a digital notch filter on the inverse Fourier transforms of the $P_i(\nu)$. The filter has negligible influence on both amplitude and line-shape of any potential axion signal, while being effective at eliminating the fluctuation artifact. In combination with the SG filter, this process ensures uncorrelated $p_{\text{proc},i}$. The $p_{\text{proc},i}$ are scaled and combined using weights that take the axion sensitivity depending on the B-field strength in individual runs into account. The resulting combined spectrum is then cross-correlated with the expected Maxwell-Boltzmann axion line-shape [117, 118] to yield the grand spectrum. A local DM velocity dispersion $\sigma_v = 156$ km/s [69] and a velocity of the laboratory relative to the DM halo $v_{\text{lab}} = 242$ km/s [119] are used, resulting in a FWHM of the expected line-shape of ~ 20 kHz.

The grand spectrum is unitless and, in the absence of an axion signal, expected from simulation to have zero mean and a standard deviation of 0.96, slightly below one due to correlations induced by the SG filter [116]. Grand spectra for both configurations are presented in Fig. 6.3. The gray dashed line shows the envelope of the expected amplitude of excesses, reflecting the SNR of an axion signal with

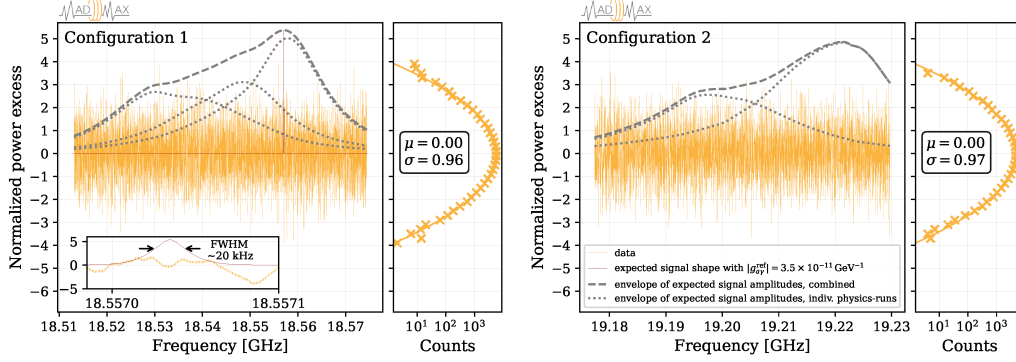


Figure 6.3: Grand spectra with 0.9 kHz bin width for both booster configurations. Dashed gray lines indicate the envelope of expected amplitude in each frequency bin for a reference axion signal with photon coupling of $|g_{a\gamma}^{\text{ref}}| = 3.5 \times 10^{-11} \text{ GeV}^{-1}$, combining all physics-runs for both configurations. Dotted lines show the expectation for individual physics runs. The expected signal shape of the same coupling strength at peak sensitivity is shown as solid line. The inset shows a zoom into this mass region. The right panels display histograms of the spectra, overlaid with a best-fit Gaussian distribution with the indicated mean and standard deviation.

$|g_{a\gamma}^{\text{ref}}| = 3.5 \times 10^{-11} \text{ GeV}^{-1}$. Its frequency dependence is determined by a combination of T_{sys} and β^2 . An axion signal would appear as a narrow localized excess with FWHM of $\sim 20 \text{ kHz}$, corresponding to 22 bins at a frequency given by m_a as shown in the inset. The largest grand spectrum excess observed in both configurations has a local significance of 4.0σ , which is consistent with statistical expectations. The probability of an equal or larger excess occurring across the entire dataset assuming no axion signal is $p = 0.08$.

Since no evidence of DM axions is found, an upper limit on the axion-photon coupling $|g_{a\gamma}|$ is derived at 95% Confidence Level (CL) for each 0.9 kHz bin as shown in Fig. 6.4. The individual uncertainty contributions on $|g_{a\gamma}|$ shown in Table 6.1 are propagated into the total uncertainty taking correlations into account. For the β^2 determination this is done using Monte Carlo simulation.

The resulting limits on DM axions exceed previous results from the CAST helioscope [56] and astrophysical considerations [120, 121].

6.5 Conclusion

Searching for axions with a dielectric haloscope, as proposed by the MADMAX collaboration, facilitates accessing the “heavy” axion mass region favored by post-inflationary PQ-symmetry breaking scenarios. A novel boost factor determination method was developed and applied to derive first results from a prototype oper-

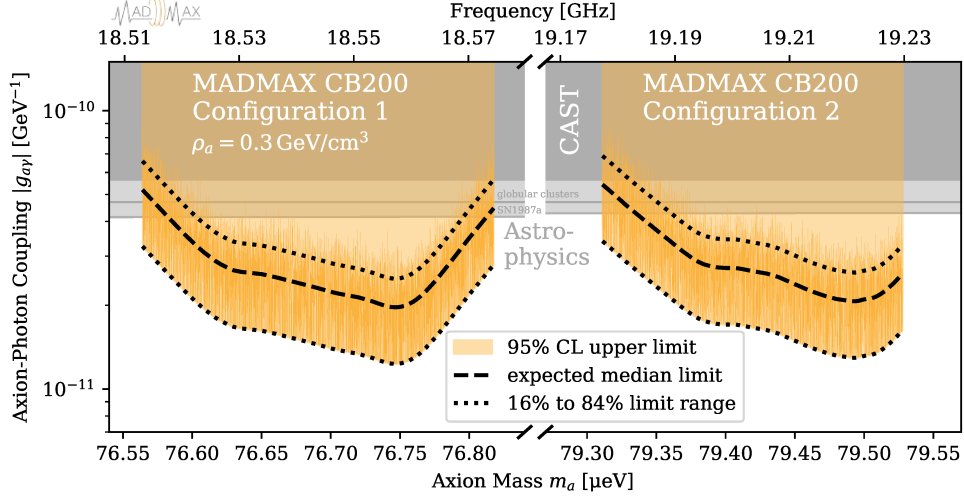


Figure 6.4: 95% CL exclusion limits in 0.9 kHz bins (orange) from the 2024 MADMAX axion search with the CB200 prototype using the Morpurgo magnet at CERN assuming a local axion DM density of $\rho_a = 0.3 \text{ GeV cm}^{-3}$. Limits are compared to the helioscope experiment CAST [56] (dark grey) as well as to supernova [120] and globular cluster limits [121] (light grey). The expected median limit and the 16% and 84% quantiles are shown as black dashed and lower and upper dotted lines, respectively. The limits are truncated below the 16% quantile.

ated in CERN’s Morpurgo magnet, using the intrinsic advantages offered by the MADMAX approach. Detailed understanding of the radio frequency response of a booster system with closed boundary conditions was demonstrated. The booster could be quickly tuned and re-calibrated to different frequencies demonstrating in principle the opportunity for future larger-range frequency scans. The data from the prototype test campaign using a small booster system inside a modest B-field allow probing for axion dark matter in previously uncharted territory. The achievements presented here and in [3] provide a firm basis for the future research and development program towards a competitive dielectric haloscope, which will include the development of boosters with more and larger dielectric disks, the demonstration of the calibration and operation of boosters at cryogenic temperatures and the implementation of tuning mechanisms via motorized disk position controls [25, 111, 85]. With ongoing development of a future, unique, large-bore 9 T dipole magnet [122, 110], the MADMAX collaboration is on a very promising track towards probing dark matter QCD axions in the 100 μeV mass region.

Table 6.1: Summary of systematic uncertainties on $|g_{a\gamma}|$. Minimum and maximum uncertainties among the frequency range for all physics-runs are reported as first and second values (when necessary). The uncertainty from each physics-run is considered individually when included in the limit estimation.

Effect	Uncertainty in $ g_{a\gamma} $
Y-factor power calibration	3 % to 5 %
Receiver chain power stability	≤ 2 %
Axion field – TE ₁₁ overlap	6 %
Boost factor determination (excl. overlap)	< 5 %
Frequency stability of TE ₁₁ mode	< 2 %
Total	5 % to 10 %

Acknowledgements

The authors would like to thank the CERN central cryogenic laboratory and the CERN magnet team for the support during all the measurements with the Morpurgo magnet. We are thankful to A. Kazempour for helping setting up the experiment at CERN and for the RFI measurements performed during data taking. We thank O. Reimann, D. Strom and A. Hambarzumjan for significant contributions during the first phase of the MADMAX project. We thank O. Rossel for providing the drawing of the setup and artistic support. We acknowledge support by the Helmholtz Association (Germany); Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) under Project No. 550641633 and Germany’s Excellence Strategy – EXC 2121 “Quantum Universe” – 390833306. Computations were performed with computing resources granted by RWTH Aachen University under project 7936. SK is supported by Fermi Research Alliance, LLC under Contract No. DE-AC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics. We acknowledge the support of the MADMAX project by the Max Planck Society.

Data availability—The data that support the findings of this letter are openly available [123]

6.6 Additional material

6.6.1 Monitoring system

A monitoring system was developed that allows to spot issues with the data early on. A screenshot at the end of the data taking of the longest run (run 1.2) is shown in figure 6.5. The overview is updated everytime a new file is recorded by the DAQ, corresponding to an interval of 15 min for the presented data taking. It tracks the measured system temperature over time as well as the physical temperature measured at different locations of the experiment. Some preliminary processing is applied to the data to show the normalized excess power and to be able to verify the expected decrease of the standard deviation of the measured noise as explained in section 4.3. The small discrepancy in integrated standard deviation between expectation and data lead to the discovery of the unexpected fluctuation discussed in section 6.4. The gaussianity plots are additional confirmation that the system noise behaves as expected. In the system temperature plot, the resonance frequency of the booster mode is marked for the average of all system temperature spectra and the last recorded spectrum to spot unexpected shifts, as the one discussed in section 5.4.2.

6.6.2 Effect of the receiver system on the boost factor

The included paper only shows the final boost factor distributions after the full procedure has been performed. In figure 6.6, the boost factor distributions as determined by the booster model only, referred to as nominal boost factor (blue dashed), are shown together with the boost factor distributions after including the effect of the receiver system (orange solid). Note that the boost factors with the effect of the receiver system are still missing the readjustment of the distance between mirror and closest disk and therefore do not match the final boost factors included in the paper. While some boost factors were marginally increased due to the receiver system, most were decreased by up to 20 %. The receiver system furthermore results in a slight shift of the frequency of maximal boost factor. The distance between receiver system and booster can in principle be adjusted, therefore for future data takings it would be sensible to spend some time on optimizing this distance for a higher boost factor. Note that the sensitivity of the experiment also depends on the system temperature, which has to be taken into account when performing such an optimization. The relative uncertainty of the boost factors changed only marginally at the frequencies of maximum boost factor. This is a result of the single parameter fit of L_{con} to the system temperature spectrum of the booster, that is shown in figure 5.15. The clear separation of maxima and minima in frequency results in a clear minimum for the optimizer and therefore small

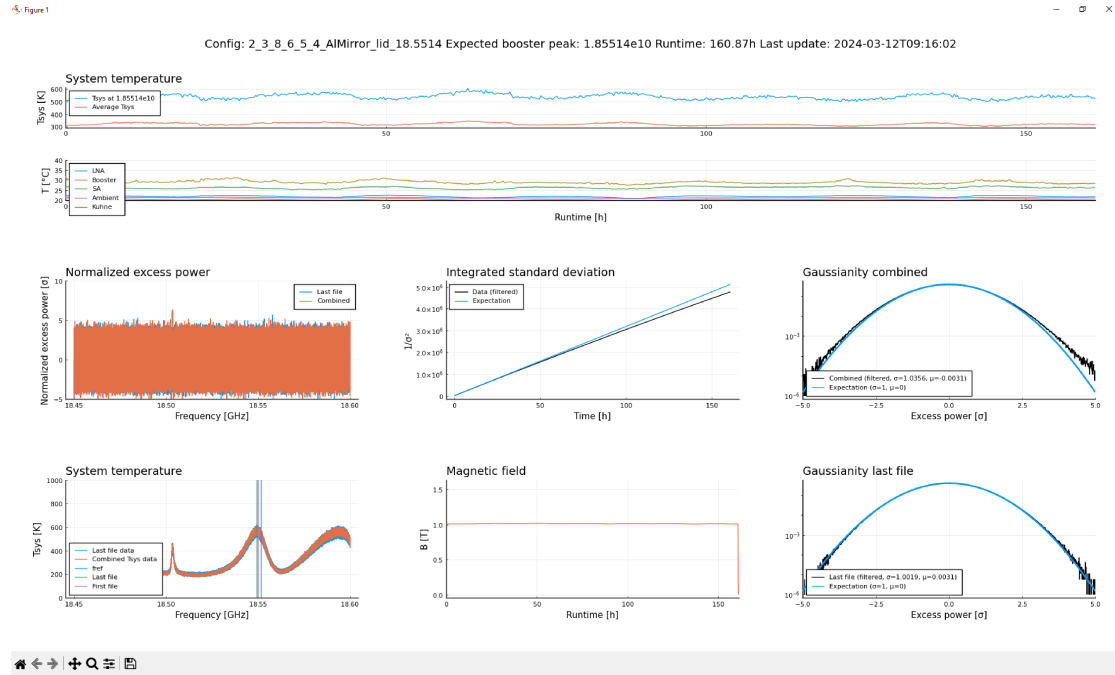


Figure 6.5: Screenshot of the monitoring system at the end of run 1.2. At the top, the application shows the average system temperature and the physical temperature over the time of the run. The panels below show the normalized excess power in units of σ , the integrated standard deviation together with the theoretical expectation, the gaussianity of the normalized excess power, the combined and last system temperature spectrum and the magnetic field.

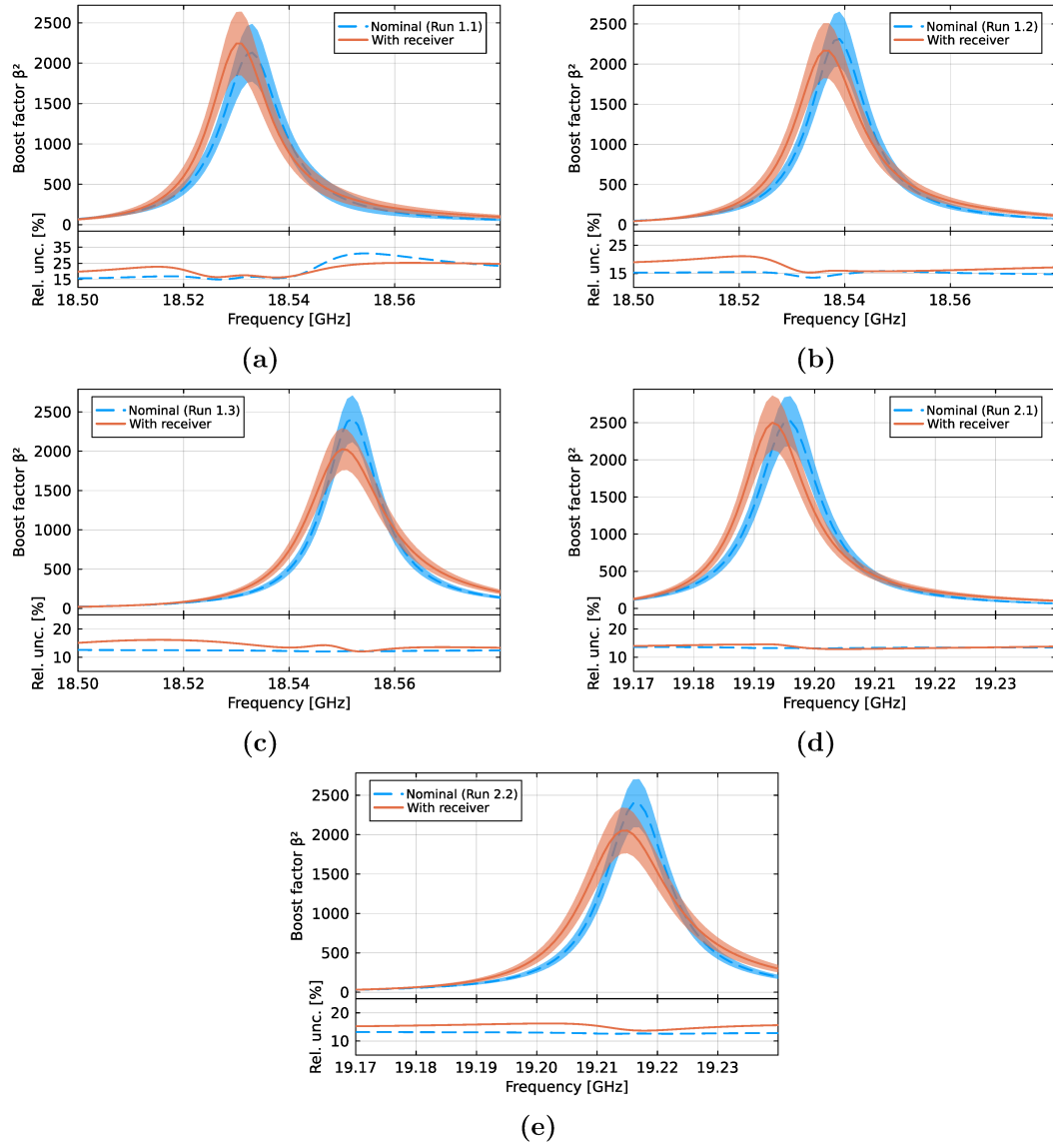


Figure 6.6: Nominal boost factor distributions (blue dashed) and boost factor distributions that include the effect of the receiver system (orange solid) for the different runs (see table 5.1) performed with CB200 in the Morpurgo magnet at CERN. The bottom part of each panel shows the relative uncertainty of the boost factor distributions, as indicated by the bands.

uncertainties on L_{con} . For the runs where no broadband measurement is available, the optimizer might choose minima for L_{con} that are separated by $\lambda/2$. Because the effect on the boost factor has the same periodicity, this does not necessarily result in an increased uncertainty.

7 Search for dark photon dark matter using an open booster

This section is published as "*First Search for Dark Photon Dark Matter with a MADMAX Prototype*" [3]. In section 7.6, additional unpublished material is presented.

Significance: A central goal of MADMAX is the development of an open dielectric haloscope. While the previous chapters focus on closed booster designs, this publication demonstrates a dark photon search using an open booster configuration. In the context of this thesis, it establishes that the developed noise model is applicable across multiple receiver and booster configurations and is therefore directly relevant for future DM searches. The determination of the receiver-induced modification to the boost factor presents the main contribution by the thesis author.

7.1 Introduction

The nature of dark matter (DM) might be the greatest unsolved mystery of particle physics and cosmology. Many experiments try to directly detect DM particles from the galactic halo in the laboratory. In recent times, very light-weight DM candidates with masses below $\sim 1 \text{ meV}/c^2$ have received increased attention [124]. One such wave-like DM candidate is the dark photon (DP), also known as hidden photon [54, 17]. DPs appear in extensions of the Standard Model (SM) that postulate an additional $U_\chi(1)$ gauge symmetry. If this symmetry is broken, the DP gains a mass m_χ . Assuming that the SM fields remain uncharged under the new $U_\chi(1)$, DPs would predominantly interact with SM particles via kinetic mixing with the ordinary photon. Its mass and small interaction strength with SM fields makes the DP a suitable DM candidate. Dark photon dark matter (DPDM) production from inflationary perturbations with $m_\chi \sim 100 \text{ peV}/c^2$ would easily saturate the observations [125, 126] and its discovery could pinpoint the scale of inflation. In the low energy limit, the interaction between DPs and photons can be described by additional source terms in the classical Maxwell's equations. In particular,

Ampère's law is modified to [53, 127]

$$\nabla \times \mathbf{H} - \dot{\mathbf{D}} = \chi \epsilon_0 (\dot{\mathbf{E}}_\chi - c^2 \nabla \times \mathbf{B}_\chi), \quad (7.1)$$

where $(\mathbf{E}_\chi, \mathbf{B}_\chi)$ are the DP electric and magnetic fields, χ is the kinetic mixing angle, and $(\mathbf{H}, \mathbf{D})$ the magnetic and displacement field of ordinary electrodynamics. For DPDM, the spatial derivative $\nabla \times \mathbf{B}_\chi$ can be neglected as the de Broglie wavelength is much larger than the setup [128, 129] whereas the temporal derivative $\dot{\mathbf{E}}_\chi$ acts as an effective current density that oscillates with frequency $\nu_\chi \approx m_\chi c^2/h$. The oscillating DP electric field drives the free charges inside a conductor which leads to emissions of ordinary photons perpendicular to its boundary. The power emitted of a metallic mirror is proportional to its area A and the local DPDM density $\rho_\chi = \frac{\epsilon_0}{2} |\mathbf{E}_\chi|^2$ [78],

$$P_0 = \chi^2 c \rho_\chi A \alpha_{\text{pol}}^2. \quad (7.2)$$

The factor α_{pol} accounts for the DP polarization and is the average fraction of DPs that the experiment is sensitive to. We assume unpolarized DPDM, setting $\alpha_{\text{pol}}^2 = 1/3$ as we are sensitive only to a single polarization [54].

In this letter, a search for DPDM using a MADMAX (MAGnetized Disk and Mirror Axion eXperiment) prototype setup is presented. MADMAX is designed to search for axions and axion-like particles in an external static B-field in the mass range of $40 \mu\text{eV}/c^2$ to $400 \mu\text{eV}/c^2$ [23, 24, 48]. With no B-field applied, the setup is still sensitive to DPs. The proposed mass range is difficult to access using traditional designs like cavity haloscopes because the effective volume where DPs (or equivalently axions) convert into photons naively scales as $\lambda^3 \propto 1/m_\chi^3$. This makes large, and therefore sensitive, conversion volumes hard to achieve. The dielectric haloscope concept of MADMAX removes the dependency of the conversion volume from the wavelength by placing dielectric disks in front of a metallic mirror [79]. This concept has already been successfully employed in searches for DPDM at lower [112] and higher [130] DP masses. The axion or DP-induced photon emission of the disks and mirror, collectively called the booster, can constructively interfere and resonate between disks. The overall increase in expected DP signal power w.r.t. a single mirror is described by the boost factor $\beta^2 = \frac{P_{\text{sig}}}{P_0}$, where P_{sig} is the signal power received by the first preamplifier of the receiver system. By controlling the spacing between disks, the boost factor can be tuned in frequency [25] allowing for both broadband and resonant searches. This enables scanning for a broad range of possible DP masses. In this work, a single fixed booster configuration at room temperature without scanning capability is used. The expected sensitivity to χ from a DP signal with bandwidth $\Delta\nu_\chi \approx 10^{-6} \nu_\chi$ can be expressed

via Dicke's radiometer equation [131] as

$$\begin{aligned} \chi = & 1.0 \times 10^{-13} \left(\frac{640}{\beta^2} \right)^{1/2} \left(\frac{707 \text{ cm}^2}{A} \right)^{1/2} \\ & \times \left(\frac{T_{\text{sys}}}{240 \text{ K}} \right)^{1/2} \left(\frac{11.7 \text{ d}}{\Delta t} \right)^{1/4} \left(\frac{\text{SNR}}{5} \right)^{1/2} \\ & \times \left(\frac{0.3 \text{ GeV/cm}^3}{\rho_\chi} \right)^{1/2} \left(\frac{\Delta\nu_\chi}{20 \text{ kHz}} \right)^{1/4}, \end{aligned} \quad (7.3)$$

where T_{sys} is the system noise temperature, Δt the effective data-taking time, and SNR is the Signal to Noise Ratio of a hypothetical DP signal. Here, DPs are considered to comprise all of the local dark matter density assuming $\rho_\chi = 0.3 \text{ GeV/cm}^3$ [10]. Both β^2 and T_{sys} are functions of frequency that depend on the booster configuration.

7.2 Experimental Setup

A schematic of the setup is shown in Fig. 7.1. Three sapphire disks with relative permittivity $\epsilon = 9.3(1)$ and thickness $1.00(2) \text{ mm}$ as well as an aluminum mirror make up the booster of the dielectric haloscope. The disks and mirror have a diameter of 30 cm each and are held parallel to each other by mechanical spacers in a fixed configuration. The planarity of the three disks has been measured with the methodology presented in [26] and gives an RMS of $50 \mu\text{m}$ ($200 \mu\text{m}$ min-max). The DP-induced emissions are coupled to a K-band Gaussian beam horn antenna via an off-axis ellipsoidal mirror of focal length $f = 504 \text{ mm}$. The horn antenna's position and orientation are precisely controlled using motorized stages. A low-noise amplifier (LNA) is connected directly to the horn antenna. Additional LNAs and bandpass filters are connected in series to further amplify the DP signal. It is then mixed down to the intermediate frequency (IF) band using a local oscillator (LO) and mixer. After a final low-pass filter, the signal is digitized with an FPGA-based DAQ using a Xilinx RFSoc4x2, analogous to the system used in [65]. The resulting IF spectrum spans from 0 to 2.4576 GHz with a resolution bandwidth of $\Delta\nu = 9.375 \text{ kHz}$. The sensitive frequency range is set by the bandwidth of the narrowest band pass filter which results in an RF range of 19.01 GHz to 20.30 GHz .

The expected DP signal power P_{sig} is proportional to the boost factor β^2 . It can be determined directly from measurement using a recently developed method [64] which relates the boost factor, primarily defined for the unknown DP-induced field, to the reflection-induced electric field E_R that is excited by a Vector Network Analyzer (VNA). This field is measured using the non-resonant bead-pull method [86] where small changes in the booster reflection coefficient Γ are related

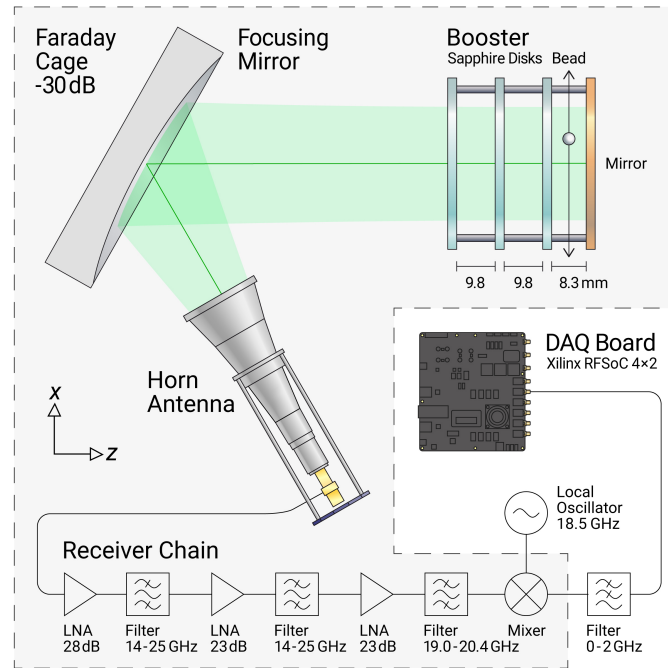


Figure 7.1: Experimental Setup. It is placed in an RF-isolated Faraday cage indicated by the dashed line. To determine the boost factor, a bead can be inserted into the booster to measure the electric field induced by a reflection measurement. Sketch not to scale.

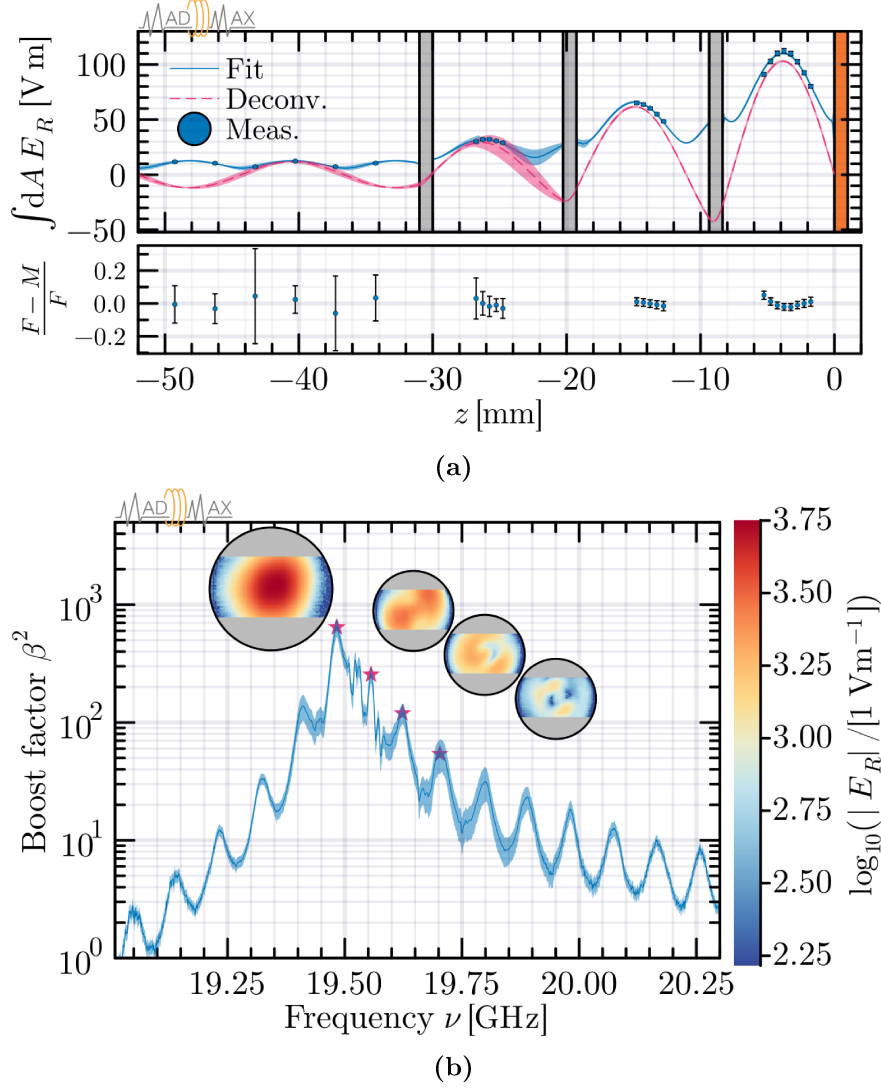


Figure 7.2: Results of the boost factor determination using the bead-pull method. **(a)** Transverse integrated electric field along the optical axis inside the booster at 19.48 GHz. Measurements (circles) fitted with a model (solid line) including the bead. The dashed line shows the field evaluated without the effect of the bead. The shaded band indicates model uncertainties from material and geometry parameters. The relative difference between fit (F) and measurement (M) is shown in the lower panel. **(b)** Boost factor as a function of frequency, including the corrections for the finite domain and receiver mismatch. The insets show the transverse electric field between the mirror and the first disk at the indicated frequencies (stars), where grey areas indicate regions unprobed due to mechanical constraints.

to the electric field at the bead's position. The general procedure to determine β^2 from bead-pull measurements is described in detail in [4, 27]. Expressed in terms of measurable quantities, the boost factor is

$$\beta^2 = \frac{P_{\text{sig}}}{P_0} = \frac{4\pi^2\epsilon_0\nu^2 F_{\text{RC}}}{8cP_{\text{in}}A} \left| \int dV E_R \right|^2, \quad (7.4)$$

where E_R is excited with input power P_{in} at frequency ν and is spatially integrated over the conversion volume. F_{RC} is a dimensionless factor that accounts for an impedance mismatched receiver system. Bead-pull measurements yield E_R relative to P_{in} such that its actual value is not required. We arbitrarily set $P_{\text{in}} = 1$ W for convenience. Figure 7.2 shows the longitudinal (a) and transverse field distribution (b) inside the booster obtained from bead-pull measurements as well as the derived boost factor as a function of frequency.

The measured field is first integrated in the transverse direction. The absolute value of the integrated electric field is then fitted separately for each frequency by a model that takes the finite size of the bead into account, shown in Fig. 7.2a. The shaded band indicates the model uncertainty from geometry and material parameters of both booster and bead. The fit allows us to obtain the deconvoluted field, i.e., the field without the response of the bead, to interpolate between measurements, and to recover the phase. The remaining integration in the longitudinal direction is then performed with the deconvoluted field, yielding β^2 via Eq. (7.4), shown in Fig. 7.2b. On the main resonance at 19.48 GHz, the transverse field is the fundamental Gaussian mode with waist radius $w_0 = 91(6)$ mm that is expected from the optical system and which has a good overlap with the uniform $\mathbf{E}_\chi$. A mix of higher-order transverse modes can also resonate inside the booster as well as between the antenna and booster, causing the additional smaller peaks, first studied in simulation in [67].

The finite domain of the bead-pull measurements does not cover the full transverse extent of E_R due to mechanical constraints from the booster leaving some of the top and bottom fringe areas of the disks unprobed (indicated by the grey areas in Fig. 7.2b). We account for these areas in the integral in Eq. (7.4) by extrapolating the electric field. This increases the naive estimate of β^2 by about 70%. The finite domain correction has been checked against independent measurements with full coverage of the field [27] and contributes an additional 10% to the systematic uncertainty of β^2 .

The receiver chain has an impedance mismatch to the antenna, resulting in an input reflection coefficient of $|\Gamma_{\text{RC}}| \sim 0.25$. Since part of the expected DP signal would be resonating between receiver chain and booster, the boost factor is further changed by the factor $F_{\text{RC}} = \frac{1-|\Gamma_{\text{RC}}|^2}{|1-\Gamma_{\text{RC}}\Gamma|^2}$. This resonance also affects the system noise temperature T_{sys} which is determined by a Y-factor calibration of

the receiver chain [80], establishing its gain and equivalent noise temperature T_e . The overall system noise temperature appears as a standing-wave pattern with a peak-to-peak system temperature of 120 K to 332 K over the measurement span of 1.2 GHz, which is dominated by the receiver chain noise of $T_e = 119(28)$ K interfering with itself. This is compatible with the expected noise of the first stage amplifier of around 120 K [132]. Only at the resonance frequency of the booster does its thermal radiation due to physical temperature significantly contribute to the system noise temperature by around 40%. The modulation of T_{sys} is visible in the received power excess in Fig. 7.3 where a higher T_{sys} leads to larger fluctuations around the baseline.

To properly quantify F_{RC} , the difference in electrical length between the measurements of Γ and Γ_{RC} needs to be known. It is extracted from fitting a 1D model to the system noise temperature, taking Γ_{RC} and Γ as input and simulating the emitted LNA noise as per [133]. The resulting value of 16.016(2) mm matches the length of the adapter used to connect the receiver chain to the antenna. Correcting for the mismatch increases the systematic uncertainty of β^2 by less than 1%. Depending on frequency, F_{RC} ranges between 0.6 and 1.6, and can consequently increase or decrease the boost factor. By tuning the distance between the antenna and booster, it is assured that on resonance frequency the boost factor is increased by $F_{\text{RC}} \approx 1.3$ to a maximum value of $\max \beta^2 = 640(110)$. The boost factor shown in Fig. 7.2b includes both corrections from the finite domain of bead-pull measurements and receiver chain mismatch. The shaded band indicates the one standard deviation uncertainty which ranges from 10 to 36%, depending on frequency.

7.3 Data Taking Run

The data-taking run lasted 16.5 days from 2023-12-22 to 2024-01-08. In this period, $N_{\text{tot}} = 9.484 \times 10^9$ spectra were measured corresponding to an effective data-taking time of $\Delta t = 11.7$ d. The spectra were averaged in batches of $N_{\text{av}} = 4.00 \times 10^6$ and saved every 10 minutes resulting in 2371 files. The setup was located in a shielded laboratory (SHELL) at DESY/University of Hamburg with coordinates: 53.58N/9.89E. The antenna is sensitive only to the zenith-pointing polarization of DPDM. The optical axis of the booster has an azimuth angle of 65° relative to north.

The data-taking routine employed the same LO hopping scheme as described in [65] to smear and suppress radio frequency interference (RFI) in the IF band. Approximately once per second (every 10^4 spectra), the LO frequency is randomly shifted within a window of ~ 9 MHz. The collected power spectra are then re-aligned in the RF band before averaging. This results in a smearing of any RFI in the IF band, effectively suppressing it without affecting signals in the RF band.

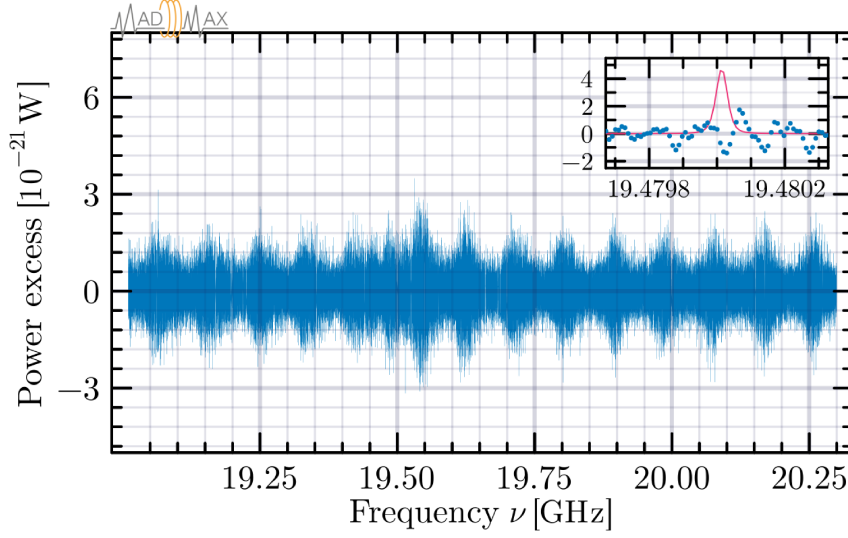


Figure 7.3: Observed cross-correlated power excess as a function of frequency. The inset shows a zoomed-in view around 19.48 GHz where the maximum boost factor occurs. A hypothetical DP signal with $\chi = 2 \times 10^{-13}$ (magenta solid line) is superimposed on observations (blue circles).

Additionally, roughly 50 bins that contain strong RFI signals in the IF band are digitally masked. The frequency hopping also causes these masks to spread in the RF band, rendering their impact on sensitivity negligible.

Communication delays between DAQ and LO are the dominant contribution to the dead time of $\sim 30\%$ and could be improved in the future. In the RF band, RFI signals are further suppressed by ~ 30 dB using an additional Faraday cage and RF absorbers. The response of these signals to additional shielding excludes a DP origin. They are removed from the analysis and are listed in the supplementary materials. A plausible source is the surface-monitoring radar of the nearby Hamburg airport.

7.4 Analysis

The analysis procedure and nomenclature closely follows HAYSTAC [116]. All saved spectra $P_i(\nu)$ are each filtered with a 6th-order Savitzky-Golay (SG) filter with a window length of $1.04 \text{ MHz} \gg \Delta\nu_\chi$ that would leave a potential DP signal intact. From this we obtain baselines $P_{bl,i}(\nu)$ and processed spectra $p_{proc,i}(\nu) = (P_i/P_{bl,i} - 1)$. For perfect baseline subtraction, the standard deviation σ_i of each processed spectrum, excluding RFI bins, is expected to be $1/\sqrt{N_{av}}$. We observe σ_i to be between 97.9% and 99.2% of this value, indicating the baseline subtraction

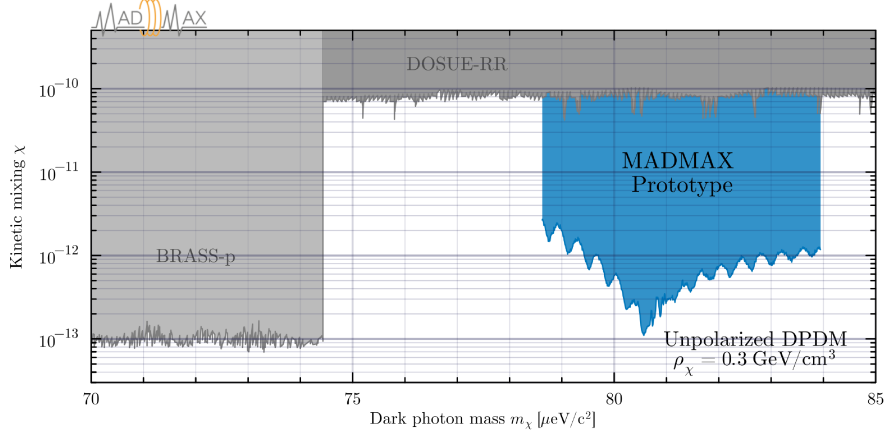


Figure 7.4: 95% CL upper limit on the dark photon kinetic mixing angle χ obtained with the MADMAX prototype as compared to the dish antenna experiments DOSUE-RR [134] and BRASS-p [105], rescaled to a common value of the assumed local dark photon dark matter density of $\rho_\chi = 0.3 \text{ GeV/cm}^3$. Unpolarized dark photons are assumed.

has a minor impact on our sensitivity. The SG filter reduces the SNR of a potential DP signal by $\eta_{SG} = 0.91$ which is obtained by running the analysis on simulated data with injected synthetic DP signals. During the run, the baseline drifted by no more than 5% in amplitude which can be attributed to gain variation caused by room temperature changes of the same order. In frequency, spectral features of $P_{bl,i}$ shifted by less than 1 MHz. Assuming the same frequency stability for β^2 , which was determined after the run, translates to a relative systematic uncertainty of $\sim 4\%$ in amplitude.

The processed spectra are combined using weights that take the expected SNR of a hypothetical single-bin DP signal into account for every bin of each spectrum [116]. The resulting combined spectrum is further cross-correlated with the expected DP line shape [118, 117] to account for the fact that the expected DP signal would stretch over ~ 5 bins. We use $\sigma_v = 154.1 \text{ km s}^{-1}$ [119] for the velocity dispersion of the local DM halo and $v_{\text{lab}} = 242.1 \text{ km s}^{-1}$ [135] for the velocity of the laboratory with respect to the local standard of rest. The unknown relative alignment between the frequency bins and DP line shape leads to variation in the discretization of the line shape, affecting the cross-correlation, and is treated as a systematic uncertainty. The observed cross-correlated power excess P_{exc} is shown in Fig. 7.3 revealing a sensitivity to DP signals of $\sim 2 \times 10^{-21} \text{ W}$. A hypothetical DP signal with $\chi = 2 \times 10^{-13}$ is shown in the inset along with a zoomed-in view of the observations around 19.48 GHz where the maximum boost factor occurs.

The largest excess has a local significance of 4.3σ . This is within the expectation

Table 7.1: Summary of systematic uncertainties. Min-to-max is shown for frequency-dependent quantities.

Effect	Uncertainty on χ
Boost factor determination	
Bead-pull measurements	2 to 17%
Bead pull finite domain correction	5%
Receiver chain impedance mismatch	<1%
Subtotal	5 to 18%
Y-factor calibration	4%
Power stability	3%
Frequency stability	2%
Line shape discretization	4%
Total	9 to 19%

of observing thermal noise only as the probability of observing a local excess at least this large is $p \approx 0.21$ for the full dataset. From the non-observation of any DP signal, we derive a 95% confidence level (CL) upper limit on $P_{\text{exc}} \rightarrow P_{\text{exc}}^{95}$ for groups of 210 neighboring bins using the approach described in [66]. The limit on the kinetic mixing angle χ is obtained by setting $P_{\text{exc}}^{95} = \eta_{SG} \beta^2 P_0$. Systematic uncertainties, summarized in Table 7.1, are assumed to be independent and propagated to χ . The most conservative value of χ within this uncertainty is adopted for the final limit.

7.5 Conclusion

The 95% CL limit on the dark photon kinetic mixing angle χ is shown in Fig. 7.4. Unpolarized dark photon dark matter ($\rho_\chi = 0.3 \text{ GeV}/\text{cm}^3$) with $\chi > 2.7 \times 10^{-12}$ can be excluded for masses between $78.62 \text{ } \mu\text{eV}/c^2$ to $83.95 \text{ } \mu\text{eV}/c^2$. Using only three disks, we have improved existing limits by up to almost three orders of magnitude with a peak sensitivity of $\chi = 1.1 \times 10^{-13}$. This demonstrates, for the first time, the feasibility of the MADMAX dielectric haloscope concept and show-cases its capability of providing a large conversion volume at high frequencies while retaining resonant enhancement. Further improvements are planned: scaling the conversion volume by adding more dielectric disks, reducing system noise temperature using a cryostat, and implementing a tunable booster via motorized disk control [111, 26]. This will vastly improve both the mass range of searches for dark photons as well as increase the sensitivity to their coupling to photons.

7.6 Additional material

7.6.1 Bin masking and frequency hopping

During the data taking, a bin masking and frequency hopping technique as described in section 4.4.2 was used. The bin masking was performed every ~ 30 min with an integration time of $\tau_m \simeq 10$ s. The resulting masklist is saved together with each recorded spectrum. The LO frequency was varied within a normal distribution of widths $N_{\text{bins}} = 1000$ every $\tau_{\text{LO}} \simeq 1$ s. The LO frequency is changed after every 10 000 spectra that are averaged in hardware, corresponding to $\tau_{\text{LO}} \simeq 1$ s. For every stored file, consisting of the average of 400×10000 spectra, a list containing the used n_{shift} values is saved. The resulting bin-by-bin statistics are shown in figure 7.5. The minimum number of samples of a single bin $N_{\text{min}} = 9.38496 \times 10^9$ deviates from the value of 9.480×10^9 , which can be explained by the fact that in some regions multiple bins had to be masked. The resulting decrease in sensitivity, which scales as $\sqrt{N}$, is $< 1\%$.

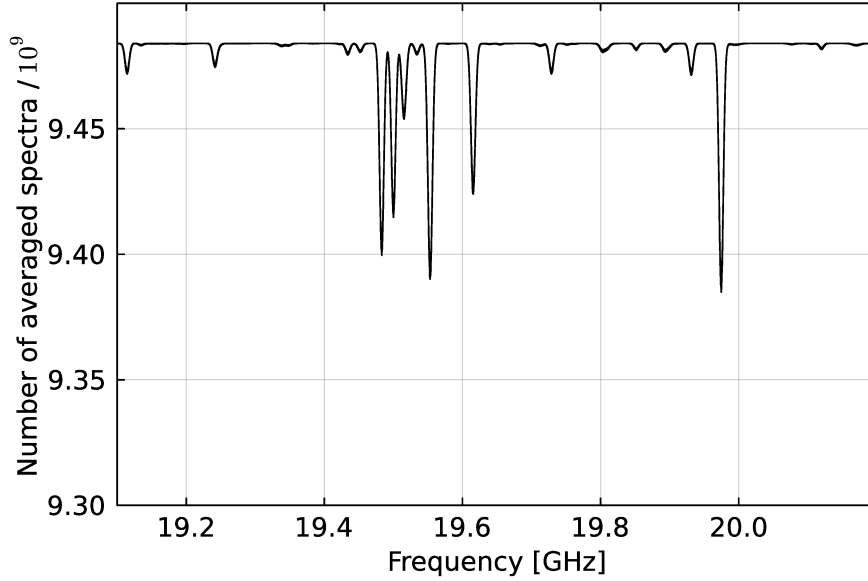


Figure 7.5: Number of points that were averaged for each frequency bin for the dark photon search. The total number of spectra measured is $N_{\text{tot}} = 9.484 \times 10^9$. Due to the bin masking, the number of averaged points was reduced to 9.38496×10^9 in the worst case around 19.97 GHz.

7.6.2 Power calibration

The power calibration of the receiver chain was performed as described for the axion search in the corresponding section 5.6.1. It results in a calibration factor

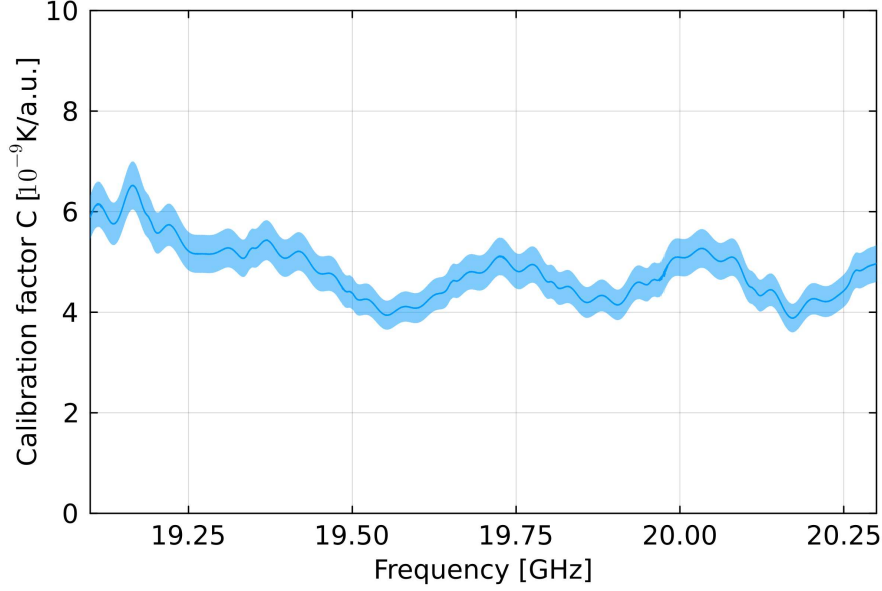


Figure 7.6: Calibration factor, converting power spectra of arbitrary unit to system temperature, including its uncertainties.

with an uncertainty of $\sim 10\%$, as shown in figure 7.6 and an equivalent noise temperature of the receiver system of $T_e = 119(28)$ K. The uncertainty is dominated by the uncertainty placed on the $ENR \simeq 14.500(75)$ dB of the noise diode and the attenuation $A = 2.9(3)$ dB of the attenuator between the noise diode and the receiver chain. To verify that the amplifiers of the receiver system do not saturate, the linearity of the measured power was verified by connecting different attenuators between the noise diode and the receiver system. In figure 7.7 the expected noise temperature T_N as per equation (5.11) is plotted against the measured noise power and compared to a linear fit. The fit and measurements match within uncertainties, confirming the operation of the components in their linear regime.

7.6.3 Application of the noise model

The determination of F_{RC} is based on the noise model already presented in section 5.2.2. Its parameters are extracted from a fit to measured power spectra of known standards, as shown in figure 7.8. The load, open and short standards

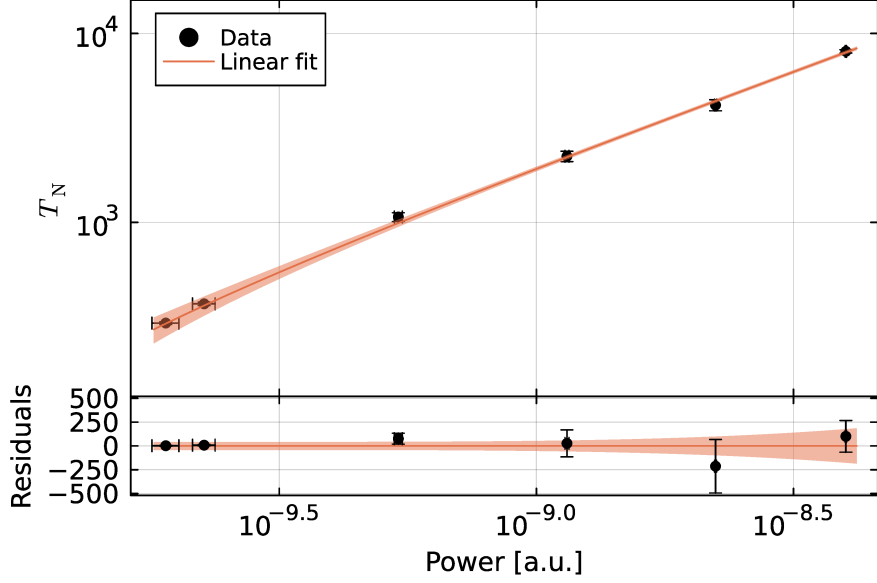


Figure 7.7: Expected equivalent noise temperature over measured power for noise measurements of different power (black dots) and the resulting linear fit (orange).

are modeled by their idealized reflectivities $\Gamma_{\text{load}} = 0$, $\Gamma_{\text{open}} = 1$ and $\Gamma_{\text{short}} = -1$, respectively. The model achieves an excellent match within uncertainties with the data. The fitted parameters are shown in table 7.2. They are of the same order of magnitude as the parameters determined for the LNA used in the axion search, shown in section 5.6.2. Only the fitted distances between the LNA and the standard are a bit shorter. They determine the distance between minima and maxima in frequency of the open and short power spectra. Since the frequency range of the power spectra (see figure 7.8) does not include multiple minima or maxima, the optimizer might not have chosen the ideal minimum for these parameters. This however does not affect the models capability to accurately describe the noise of the LNA within the displayed frequency range. The model is then used to re-

Table 7.2: Noise model parameters extracted from the fit of the model to measurements of known standards.

V_n [pV]	I_n [pA]	$ c $	ϕ_c	L_{short} [ps]	L_{open} [ps]
580(25)	7.0(4)	0.6(1)	314(10)	5.40(1)	29.1(1)

produce the system temperature spectrum of the booster, as shown in figure 7.9, by modeling the booster by its measured reflectivity Γ . The model reproduces the data within uncertainties over almost the full frequency range, including the slightly skewed shape of the broadband oscillation. Only around the booster res-

onance at ~ 19.5 GHz, the data's amplitude is slightly underestimated. This does however not influence the fit of the distance between LNA and booster L_{con} . The parameter only affects the distance between minima and maxima in frequency, which is accurately reproduced. From the fit, $L_{\text{con}} = 53.423(6)$ ps $\simeq 16$ mm is determined, which matches the physical length of the connector used between the two devices. The resulting correction factor on the boost factor F_{RC} , as defined in section 7.2, oscillates between ~ 0.75 and ~ 1.5 . This shows that the effect can potentially be exploited to increase the sensitivity. In the present case, the boost factor maximum was increased by ~ 25 %.

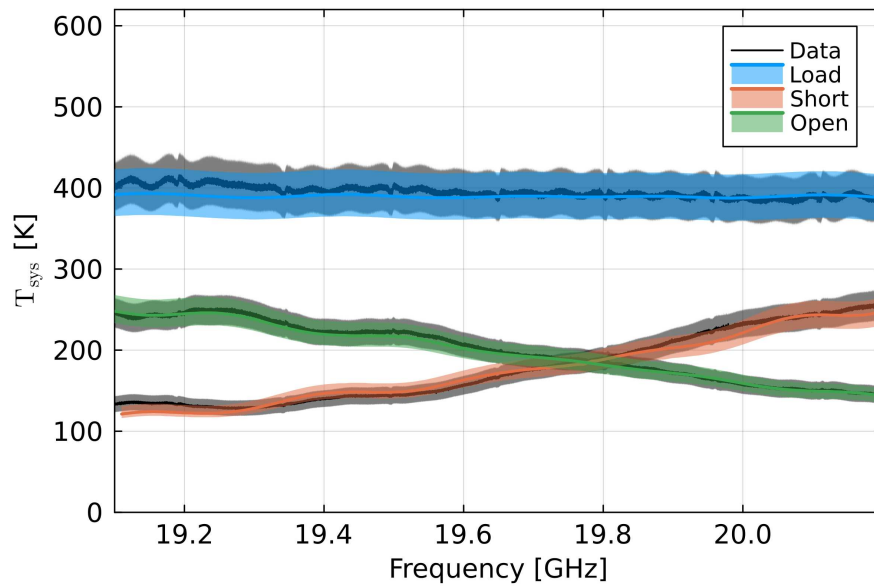


Figure 7.8: Measured and modeled noise spectra T_{sys} of the receiver system connected to open, short and matched load standards. The bands around the curves show their associated uncertainty, stemming from the power calibration of the data. The load (blue) is close to constant over the frequency range, while the short (orange) and open (green) measurements exhibit an oscillation of opposite phase. The simulation matches the measurements within uncertainties.

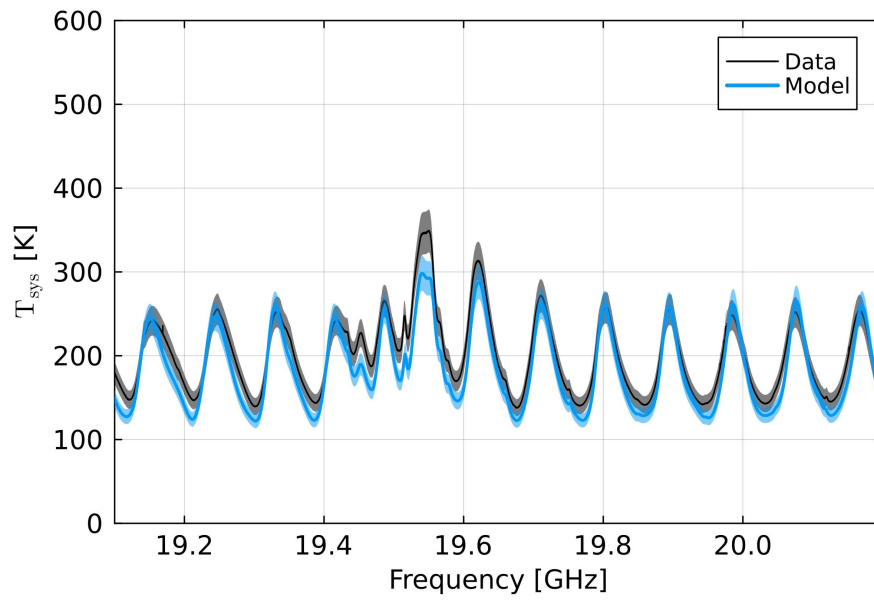


Figure 7.9: Measured and modeled noise spectrum of the OB300 booster connected to the receiver system. The model (blue) reproduces the shape of the oscillation that is visible in the data (black) and slightly underestimates the amplitude around the booster resonance at 19.5 GHz.

8 Conclusion

Dielectric haloscopes represent a new experimental approach to detecting axions in the mass range of $\sim 40\,\mu\text{eV}$ to $\sim 400\,\mu\text{eV}$. This range is well-supported by theory but has remained largely unexplored to date. Determining the sensitivity of a dielectric haloscope to dark matter, as quantified by its boost factor, remains a significant challenge. Full-wave simulations are no longer able to reproduce the exact electromagnetic response of such large-scale systems and can take weeks to perform even for simplified setups. Furthermore, the receiver system necessary for a dark matter search alters the boost factor by reflecting part of the potential signal. Therefore, the receiver system is not merely a passive component, but actively modifies the electromagnetic response of the haloscope.

In this thesis, two complementary models are developed which, when combined, enable the determination of the boost factor of a closed dielectric haloscope in the presence of a mismatched receiver system.

The first is a single-mode model of a closed booster, which is capable of reproducing the measured reflectivity of the booster and determining the corresponding boost factor. This requires fitting its effective parameters to account for three-dimensional effects, such as mirror tilt and disk deformation. Validation against full-wave simulations of a three-disk booster demonstrates that mirror tilts up to $0.05\,\text{deg}$ and realistic disk deformations with a min-max flatness of $\sim 50\,\mu\text{m}$ are described within an uncertainty of $\sim 12\%$ on the boost factor.

The second model describes the noise emitted by the receiver system and accurately reproduces measured power spectra of the booster by fitting the electrical distance between the two systems. From the extracted parameter, the effect of the receiver system on the boost factor can be quantified. This highlights that the boost factor cannot be treated purely as a property of the optical system, but also depends on the connected receiver. The noise model is further shown to be able to relate unexpected changes of the system temperature during a data taking to a specific model parameter and quantify the resulting change in boost factor.

These models were applied to a MADMAX prototype called CB200 that was taken to CERN to perform an axion search in the 1.6 T Morpurgo magnet. The spacings of the three $\varnothing 200\,\text{mm}$ sapphire disks could be reconfigured by exchanging the separation rings between them, which was used to perform axion searches around 18.55 GHz and 19.21 GHz. By applying pressure against the mirror using a tuning rod, its frequency of sensitivity was further adjusted by $\mathcal{O}(10\,\text{MHz})$. This

resulted in three runs around 18.55 GHz and two around 19.21 GHz, demonstrating for the first time a limited tuning capability of a dielectric haloscope. By applying the two models outlined above, the corresponding boost factors were determined. The receiver system was found to reduce the boost factors by up to 20 %. In the absence of an axion signal, limits on the axion parameter space were placed. The limits down to an axion photon coupling of $g_{a\gamma} \simeq 2 \times 10^{-11} \text{ GeV}^{-1}$ in the axion mass range 76.56 μeV to 76.82 μeV and 79.31 μeV to 79.53 μeV already surpass previous limits in the same mass range.

In the future, MADMAX plans to use an open booster to allow for easier movement of the disks and thus tune the system for different axion masses. The booster model presented in this work does not apply to an open booster, because additional higher-order modes occur that are mainly induced by antenna reflections. To determine the boost factor, the measurement based bead-pull method was developed and applied to an open booster with three $\varnothing 300 \text{ mm}$ sapphire disks in the PhD thesis by J. Egge. This thesis contributed to the subsequent dark photon search by configuring and characterizing the receiver system. To avoid signals in the IF band from interfering with the data taking, a bin masking and frequency hopping technique was successfully implemented. By applying the noise model, the maximum boost factor was found to have increased by about 25 % due to the mismatched receiver system. The resulting limits on the dark photon mass range 78.62 μeV to 83.95 μeV with a peak sensitivity to the kinetic mixing of the dark photon with the SM photon of $\chi \simeq 2 \times 10^{-13}$ already surpass previous limits by up to three orders of magnitude.

The two methods of determining the boost factor of a dielectric haloscope provide individual tradeoffs: The measurement based bead-pull method, developed by J. Egge, enables the accurate characterization of higher-order mode contributions. However, it takes at least several days to perform for a three disk setup and is mechanically challenging. By increasing the number of disks and going to cryogenic environments, the needed time will further increase and the mechanics of moving a bead through the setup will become more complicated. The model based method developed as part of this thesis can be performed in a matter of minutes, but only considers the contribution of a single mode to the boost factor. It was furthermore only validated against full-wave simulations of a small-scale booster with three $\varnothing 10 \text{ cm}$ disks. An increase of the booster radius is not expected to affect those results, but stronger resonances due to a higher number of disks can in principle lead to stronger field deformations. They also lead to a deeper dip in reflected power $|\Gamma|^2$. In the studied case, the difference between the expected dip in $|\Gamma|^2$ due to dielectric losses and the dip in the simulated reflectivity of the strongest mirror tilt that still led to an accurate prediction of the boost factor was $\sim 4 \text{ dB}$. This deviation defines a practical limit beyond which the model might

no longer produce the correct boost factor and further validations are necessary. This does not generally limit the usage of more disks – by making sure everything is properly aligned one might still be able to reach stronger resonances without exceeding the limits of the model. Alternatively, more disks can also be used to achieve a broader boost factor rather than increasing the resonance.

The presented noise model is not only interesting for dielectric haloscopes but is broadly applicable to experiments that rely on mismatched receiver systems, commonly encountered in microwave-based DM searches. Across both searches, the receiver system was found to modify the boost factor by $\mathcal{O}(10\%)$, demonstrating that this is a significant effect. In the case of the dark photon run, the boost factor maximum was increased by about 25%, suggesting that it could be deliberately exploited to enhance the experimental sensitivity. Furthermore, the demonstrated capability to extract model parameters from the system temperature without interrupting the data taking potentially allows for longer integration times by avoiding frequent calibration measurements.

While clearly demonstrating the viability of dielectric haloscopes, the presented axion limit is still about three orders of magnitude away from the envisioned sensitivity to the QCD axion. MADMAX is actively working on performing experiments at cryogenic temperatures, which is expected to increase the experiment's sensitivity by roughly an order of magnitude. The collaboration is furthermore working on an intermediate magnet with a 4 T magnetic field. Together with an increased disk radius of 15 cm, this leads to an increase in sensitivity by another order of magnitude. With a longer measurement time of ~ 30 d and an increase in boost factor to $\sim 20\,000$ this allows for reaching sensitivity to the KSVZ axion. To enable such a boost factor, it is planned to extend OB300 to 20 disks. Combining the presented booster model with bead-pull measurements provides a promising path towards determining the sensitivity of such large-scale setups. The model can in principle be extended to include a higher number of modes, while reduced bead-pull measurements could be used to verify the existence of only those modes. Alternatively, a simplified model could be used to propagate the field measured at a single plane in front of the booster into the booster structure.

In summary, this thesis demonstrates that dielectric haloscopes provide a viable method for probing the axion and dark photon parameter spaces. The developed models enable the accurate determination of the boost factor of closed booster systems in the presence of a mismatched receiver. The results show that the boost factor is not purely a feature of the booster, but must be understood as a property of the coupled system consisting of booster and receiver. This insight together with the developed methods provides a solid foundation for future DM searches by the MADMAX collaboration.

List of own contributions

This cumulative thesis consists of three core publications, listed below together with the main contributions by the thesis author ("*Own contributions*").

Sensitivity of a closed dielectric haloscope to axion dark matter [1]

Own contributions:

- Development of booster and receiver noise model
- Development of the procedure to determine the boost factor and its uncertainties
- Validation against simulations provided by a collaborator
- Application of the procedure to data taken with the CB200 MADMAX prototype
- Leading role in paper writing

First Search for Axion Dark Matter with a MADMAX Prototype [2]

Own contributions:

- Determination of the boost factor distributions
- Co-development of the analysis framework
- Implementation of the monitoring system
- Support of the data taking
- Significant contributions to the paper writing

First Search for Dark Photon Dark Matter with a MADMAX Prototype [3]

Own contributions:

- Determination of the receiver system's effect on the boost factor
- Setup of the data acquisition system
- Implementation of frequency hopping and bin masking
- Support of the data taking
- Contributions to the paper writing

Additional co-authorships with minor contributions**Experimental determination of axion signal power of dish antennas and dielectric haloscopes using the reciprocity approach [4]**

Own contribution: Limited involvement in performing the required measurements

A new Limit for Axion Dark Matter with SPACE [5]

Own contribution: Support in implementing the receiver system and DAQ

Constellation: The autonomous control and data acquisition system for dynamic experimental setups [6]

Own contribution: Implementation of Constellation as the slow control system for the MADMAX cryostat

Axion Dark Matter EXperiment 2A Cavity Characterization [7]

Own contribution: Support in implementing the beadpull method on the ADMX cavity.

Acknowledgments

I thank my supervisors Erika Garutti and Axel Lindner for their continuous support and the trust they put in me and my work. Thanks to the MADMAX collaboration for the many years we now spent together and the great people I got to know. A special thanks goes to Jacob Egge, who has been a constant source of scientific advice since his supervision of my master thesis.

Besides the scientific part, I am grateful for all the mental support I received – especially from my partner Elisabeth, who had to watch me go on *Physikerausflüge* having a good time and still deal with my bad mood at home when the work was more stressfull. Of course, I am also very happy to have an amazing group of friends and a great family. Special shoutout to the *Weihnachtsmarkt* and *Bergziegen* 😊

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