

# CuMnAs/Al and MnTe/Al Magnetic Material-Superconductor Heterostructures

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# Abstract

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Magnetic material-superconductor heterostructures are promising candidates for applications in superconducting electronics and superconducting spintronics. This thesis investigates bar devices of two different magnetic material-superconductor heterostructures: CuMnAs/Al and MnTe/Al, where CuMnAs is an antiferromagnet (AFM) and MnTe is an altermagnet. AFMs and altermagnets are interesting for use in spintronics due to their net zero magnetization resulting in no stray fields. During measurements the MnTe/Al heterostructure showed directional dependence of the sheet resistance, and logarithmically increasing resistance with decreasing temperature. No superconductivity was observed, which is speculated to be from magnetic moments in MnTe suppressing superconductivity. The temperature dependence of the resistance could be due to weak localization or the Kondo effect. The CuMnAs/Al sample showed strong hysteresis with regards to field history. An out of plane field was used to non-monotonically shift the peak switching current as a function of in plane field, and a W-shaped switching current as a function of in-plane field was found for both studied devices depending on the field history. Furthermore, a large zero-field superconducting diode effect (SDE) ( $\eta \geq 18\%$ ) for both studied devices was found, with a maximum of  $\eta = 22.7\%$ . This result makes the CuMnAs/Al heterostructure a promising candidate as a zero-field superconducting diode without stray field. The origin of the zero-field SDE as well as of the out of plane field history effects on the switching current behavior is speculated to be due to uncompensated magnetic moments near the CuMnAs/Al interface. The W-shaped switching current as a function of magnetic field might be explained by spin canting.

# Populärvetenskaplig sammanfattning

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## Samspel mellan magneter och supraleutare

**Magneter och supraleutare kan kombineras för att skapa nya material med lovande egenskaper för applikationer inom kvantteknologi. I avhandlingen undersöks två sådana material, varav ett visar sig fungera som en supraledande diod.**

Under de senaste åren har kvantteknologi vuxit fram som ett nytt område där det satsas stora pengar inom Sverige och resten av världen. Detta är teknologi där kvantfysikens effekter dras nytta av på olika sätt, exempelvis genom kvantdatorer, väldigt känsliga magnetmätare, eller elektronik som kan överföra information utan energiförluster. Många av dessa tekniker är baserade på så kallade supraleutare, vilka är spännande eftersom de uppvisar kvantmekaniskt beteende även i relativt stora system. Dessa effekter visar sig i supraleutarnas oändligt stora ledningsförmåga samt genom att det inte kan finnas något magnetfält i deras innandömen. Supraleddning är en egenskap som vissa material, exempelvis aluminium, får när de kyls ned till väldigt låga temperaturer. Ofta krävs det att de kyls ned till några enstaka kelvin, där 0 K är den absoluta nollpunkten ( $-273.15^{\circ}\text{C}$ ). Detta är väldigt låga temperaturer, och mätningarna i den här avhandlingen har gjorts i en speciell kyl som kan nå temperaturer på bara några tusendelars grad över den absoluta nollpunkten.

Traditionellt har supraleutare slutat uppvisa kvantegenskaper när de läggs i närheten av magneter. På senaste tiden har dock samspelet mellan supraleutare och magneter undersökts vidare, och visats leda till nya exotiska typer av supraleddning som kan vara av nytta inom spinntronik: teknik som utnyttjar elektroners spinn för att överföra eller lagra information. Det har också visat sig att magneter kan användas för att bygga supraledande dioder genom att läggas i nära anslutning till supraleutarna. Dessa dioder kan sägas vara den supraledande motsvarigheten till vanliga dioder,

och karakteriseras av att supraledarens oändliga ledningsförmåga endast gäller i en riktning. Sådana dioder är viktiga byggstenar för utvecklingen av framtidens elektronik.

Hittills har forskningen på samspelet mellan supraledare och magneter mest fokuserat på så kallade ferromagneter. Dessa är de magneter man stöter på i vardagen, som exempelvis kylskåpsmagneter. I det här arbetet har istället ett antiferromagnetiskt och ett altermagnetiskt material utforskats. Antiferromagnetiska och altermagnetiska material har inte varit kända lika länge som ferromagneter, eftersom de inte avger några magnetiska fält. Trots detta anses de vara magnetiska material eftersom det finns en slags ordning i hur elektronernas spinn ligger i materialen. I icke-magnetiska material ligger de istället huller om buller. Altermagnetiska och antiferromagnetiska material skiljer sig endast på ett väldigt subtilt sätt från varandra, vilket är varför altermagnetism inte upptäckts förrän nyligen. Eftersom dessa material inte avger några magnetiska fält, påverkar de inte andra komponenter som kan vara känsliga för magnetfält, vilket gör att de är lovande för användning inom spintronik.

I denna avhandling har två olika material undersökts. Ett som kombinerar en antiferromagnet med aluminium, och ett som kombinerar en altermagnet med aluminium. Det undersökta materialet med altermagneten visade ingen supraleddning, vilket kan betyda att altermagnetismen förstörde supraleddningen. Materialet med antiferromagneten visar oväntade egenskaper där supraleddningen förstärks vid större yttre magnetiska fält. Det visar sig också vara en potentiell kandidat för användning som supraleddande diod, med stor diodeffekt.

# Preface

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This thesis was conducted at the Center for Quantum Devices at the University of Copenhagen during the spring semester of 2026. I would like to start by thanking my supervisor Saulius Vaitiekėnas as well as Lucas Casparis for their encouragement, support, and engagement during my stay at QDev. Their enthusiasm has been infectious, and without their help this thesis would not be possible. I also want to thank August Jacobi for all the discussions about what to measure next, for the many hours spent by the fridge, and for (somewhat successfully) teaching me how to make latte art. A big thank you to everyone else at QDev for making this semester so enjoyable and for making me feel at home. Also thanks to my supervisor from Lund, Martin Leijnse, for continuous support during my thesis. Finally, I would like to thank all of my friends and family who have supported me during these five years of studying in Lund.

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# Introduction

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Superconductivity has, since its discovery in 1911, come to be used in several applications, from magnetic resonance imaging (MRI) and nuclear magnetic resonance (NMR) [1] to quantum technologies such as qubits and superconducting quantum interference devices (SQUIDs) [2, 3]. The usefulness of superconductors lies in their ability to conduct electricity with zero resistance and in their ability to expel magnetic fields. These properties hold up to a critical magnetic field and current, at which point the superconductivity breaks down and the material becomes normal again. The superconductors' dissipationless current transport together with improvements in cryogenic cooling efficiency makes superconducting electronics a promising alternative for efficient computing, emerging as a possible solution to the increased power consumption which has come with the recent expansion of large-scale cloud computing and big data [4, 5]. However, the usefulness of superconducting electronics in high-end computing applications relies on the optimizations of different superconducting analogs to semiconducting devices [5]. One important such analog on the road to useful dissipationless superconducting technologies is the superconducting diode, characterized by having a critical current that is larger in one direction, enabling zero resistance in only one direction [5–8]. Optimizing these superconducting diodes is an active research field, and there are several ways of realizing a superconducting diode, including applying external magnetic fields [6, 8], by creating ferromagnetic-superconductor heterostructures [8, 9], or by creating altermagnetic-superconductor heterostructures [10]. The magnetic material-superconductor heterostructure is especially attractive because of its potential to realize a superconducting diode without an external field, making it more useful in practical applications [9].

Another area in which superconductors show promise is the field of superconducting spintronics [11]. Spintronics is the branch of technology that uses the magnetic moments of electrons in addition to their charge in order to store and transfer infor-

mation [12], and superconductors have the potential of minimizing heating and dissipation in spintronic devices [11]. Magnetic material-superconductor heterostructures are of central importance in this emerging area [11], and interfaces between ferromagnets and superconductors have been used in order to generate spin triplet Cooper pairs that can carry spin information without dissipation [13].

In this thesis, the magnetic material-superconductor heterostructures MnTe/Al and CuMnAs/Al are examined, where MnTe is an altermagnet [14] and CuMnAs is an antiferromagnet [15]. Antiferromagnets in general are interesting for use in spintronics due to their robustness against external fields, their absence of stray magnetic fields, and their ultrafast dynamics [12]. CuMnAs in particular is an antiferromagnet that has gained attention due to its favorable properties for spintronics [16, 17] such as the ability to switch its magnetic moments using current pulses [18], as well as due to its relative ease of growth compared to other antiferromagnets with similar properties [19]. Altermagnets have, in addition to the absence of stray magnetic fields, a spin-split band structure, making them promising for use in spintronics [20]. MnTe is a particularly interesting altermagnet since it is a semiconductor and because it allows both p- and n-doping for hole or electron transport [21].

The motivation for this thesis is to investigate the properties of CuMnAs and MnTe in proximity with Al in order to study the interplay between magnetism and superconductivity, and to investigate whether these heterostructures might be useful for superconducting electronics/spintronics. To the best of our knowledge, these specific heterostructures had not been investigated prior to this work, so this thesis provides a foundation for further studies into these material stacks.

The thesis starts with a background in the necessary theory needed to understand the rest of the work, covering basic concepts in superconductivity and magnetism as well as providing some background on CuMnAs and MnTe. Afterwards, the measurement setup and results for the MnTe/Al heterostructure are described. After that, the growth and fabrication, along with the measurement setup used during the investigation of the CuMnAs/Al devices, are described. Following that, measurements of the switching currents are presented and commented. After the investigation into the switching currents, measurements of the zero-field superconducting diode efficiency are presented. The thesis is concluded with a summary of the results and an outlook to future possible research.

# Background

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## II.1 Superconductivity

Superconductivity is a property of certain materials which they attain via a phase transition when cooled down below a critical temperature,  $T_c$ . This state is characterized by two key properties: zero resistance to electrical current, and the Meissner effect, i.e. perfect diamagnetism and the expulsion of magnetic fields from the material when cooled down below  $T_c$  [2]. The microscopic explanation for superconductivity was presented in by 1967 Bardeen, Cooper, and Schrieffer (BCS) [22]. In their model, electrons of opposite momenta and spin:  $(\mathbf{k} \uparrow, -\mathbf{k} \downarrow)$ , pair up and form Cooper pairs. The origin of this attractive force is phonon-mediated: an electron traveling through the lattice attracts positive ions, which in turn attracts a second electron. Since Cooper pairs in ordinary BCS theory are formed from two electrons with opposite spin, they have spin 0 and are therefore composite bosons. Because they are bosons, they don't have to follow the Pauli principle, and form a many-particle condensate wavefunction analogously to a Bose-Einstein condensate (BEC). This condensate maintains phase coherence over macroscopic distances, and is the microscopic reason behind the phenomena of superconductivity [2].

The theory of superconductivity is a vast field, so this background will focus on the parts specifically relevant for this thesis. The background will start by introducing the London equations and use them to show that there must be a critical field  $H_c$  as well as the critical current  $I_c$ . The mechanism by which the critical field and current are likely reached in thin film aluminum, which is the superconductor used in this thesis, will then briefly be discussed. Finally the superconducting diode effect (SDE) will be introduced and discussed.

### II.1.1 The London Equations and Critical Current/Field

The motivation for the London equation given here follows Tinkham [2], though was originally given by F.London [23]. By using the canonical momentum:

$$\mathbf{p} = m\mathbf{v} + e\mathbf{A}, \quad (\text{II.1})$$

where  $\mathbf{A}$  is the magnetic vector potential, and noting that the ground state in the absence of an applied field should have zero net canonical momentum (as shown in a theorem by Bloch), the local average velocity can be expressed as

$$\langle \mathbf{v} \rangle = -\frac{e\mathbf{A}}{m}. \quad (\text{II.2})$$

This holds given that the superconducting electrons retains its ground state property of  $\langle \mathbf{p} \rangle = 0$ , which was postulated by London. Denoting the density of electrons in this ground state by  $n_s$  leads to the expression for the current density:

$$\mathbf{J} = n_s e \langle \mathbf{v}_s \rangle = -\frac{n_s e^2 \mathbf{A}}{m}, \quad (\text{II.3})$$

from which the original London equations can be recovered:

$$\mathbf{E} = \frac{\partial}{\partial t} \left( \frac{m}{n_s e^2} \mathbf{J} \right) \quad (\text{II.4})$$

$$\mathbf{B} = -\nabla \times \left( \frac{m}{n_s e^2} \mathbf{J} \right) \quad (\text{II.5})$$

by taking the time derivative and curl of Eq. (II.3) respectively.

From the London equations the interaction between a superconductor and a magnetic field can be derived. From Maxwell's laws in the static case, it is known that  $\nabla \times \mathbf{B} = \mu_0 \mathbf{J}$ . Taking the curl of both sides and simplifying:

$$\nabla \times (\nabla \times \mathbf{B}) = \nabla \times \mu_0 \mathbf{J} \implies \quad (\text{II.6})$$

$$\nabla(\nabla \cdot \mathbf{B}) - \nabla^2 \mathbf{B} = \nabla \times \mu_0 \mathbf{J} \implies \quad (\text{II.7})$$

$$-\nabla^2 \mathbf{B} = \nabla \times \mu_0 \mathbf{J} \implies \quad (\text{II.8})$$

$$\nabla^2 \mathbf{B} = \mu_0 \frac{e^2 n_s}{m} \mathbf{B}, \quad (\text{II.9})$$

where the following identities were used, in order: the vector identity  $\nabla \times (\nabla \times \mathbf{B}) = \nabla(\nabla \cdot \mathbf{B}) - \nabla^2 \mathbf{B}$ ,  $\nabla \cdot \mathbf{B} = 0$  from Maxwell's equations, and Eq. (II.5). Following Snoke [24], we make the ansatz  $\mathbf{B}(\mathbf{x}) = \mathbf{B}_0 e^{-z/\lambda_L}$  and insert into Eq. (II.9), arriving

at the expression for the London penetration depth,  $\lambda_L$ :

$$\frac{1}{\lambda_L^2} = \frac{\mu_0 e^2 n_s}{m}. \quad (\text{II.10})$$

This implies the Meissner effect: that superconductors expel magnetic fields over distances greater than  $\lambda_L$ .

Consider now a superconductor in an external field. As discussed, fields inside the superconductors are expelled, ensuring that the field inside the superconductor stays at zero. The external field is compensated for by surface circulating currents in the superconductor, which have an associated energy cost [24]. Note that this does not mean that there is dissipation in the system: once the compensating currents exist there is no energy required to maintain them. The energy cost occurs when creating them. At some applied field strength, then, the energy cost associated with creating the surface screening currents gets larger than the energy gain from forming the superconducting condensate, and the material turns normal. This field is called the critical field,  $H_c$ . Given that there exists a critical field  $H_c$ , one can also deduce that there must be a critical current density  $J_c$  at which the superconductor turns normal. This is because, from basic electromagnetism, a current generates a magnetic field, and if this magnetic field is higher than  $H_c$ , superconductivity is lost. As an example, consider the magnetic field generated from a current flowing through a wire, at the surface of the wire [25]:

$$B = \frac{\mu_0 I}{2\pi r}. \quad (\text{II.11})$$

Following the previous argument, this means  $I_c = \frac{2\pi r B_c}{\mu_0}$ . The current flow is limited to a depth of  $\lambda_L$ , so flows through an area of  $2\pi r \lambda_L$ , giving the critical current density:

$$J_c = I_c / (2\pi r \lambda_L) = \frac{B_c}{\mu_0 \lambda_L}. \quad (\text{II.12})$$

This is a parameter intrinsic to the material and does not depend on the geometry of the superconductor [24]. At the limit  $T = 0$  K,  $n_s$  can be taken to be the density of conduction electrons  $n$ , so it is in this limit that these derivations hold [2]. For  $T \rightarrow T_c$  it is expected that  $n_s \rightarrow 0$ , at which point  $\lambda_L$  will diverge, see Eq. (II.10).

It should be noted that this mechanism by which the critical current can be reached is only applicable for bulk type-I superconductors [26], and was described in order to argue that there must be a critical current and critical magnetic field for a superconductor. For type II superconductors, there are two critical fields:  $H_{c1}$  and  $H_{c2}$ . When the external field exceeds  $H_{c1}$ , so called Abrikosov vortices form in the material. These vortices each allow penetration of one flux quantum,  $\Phi_0 = hc/2e$ , and take the form of supercurrent flowing around a non-superconducting center

[2]. These vortices are pinned to specific locations in the material, resulting from spatial inhomogeneities in the material. When applying a current through such a superconductor, there is a Lorentz force on the vortices, and the critical current is reached when this force becomes too large for the pinning centers to hold the vortices [2]. Since thin-film Al behaves as a type-II superconductor [27], this is the mechanism ascribed to the critical currents in this thesis.

### II.1.2 Superconducting Diode Effect

As discussed, a defining characteristic of a superconductor is zero resistance. From Ohm's law  $V = IR$ , this means that as long as the sample is superconducting, there is no voltage drop over the sample. The maximal current that can be supported by the superconductor is called the critical current, as discussed in Section II.1.1, and is a property intrinsic to the material studied. During measurements, the current at which the material turns from superconducting to dissipative, or normal, is called the switching current. The switching current differs from the critical current in that it has an element of randomness due to thermal and quantum fluctuations, and in general follows some distribution dependent on sweep rate, material properties, and the geometry of the studied material [28]. A typical I-V-curve of a superconductor is illustrated in Fig. II.1, where the switching currents  $I_c^+$  and  $I_c^-$  as well as the retrapping currents  $I_r^+$  and  $I_r^-$  are shown. The retrapping current is the current at which the sample turns superconducting when sweeping from the dissipative state, and is usually smaller in magnitude than the switching currents. This is often attributed to the sample heating up in the normal state, thus degrading the superconductivity.

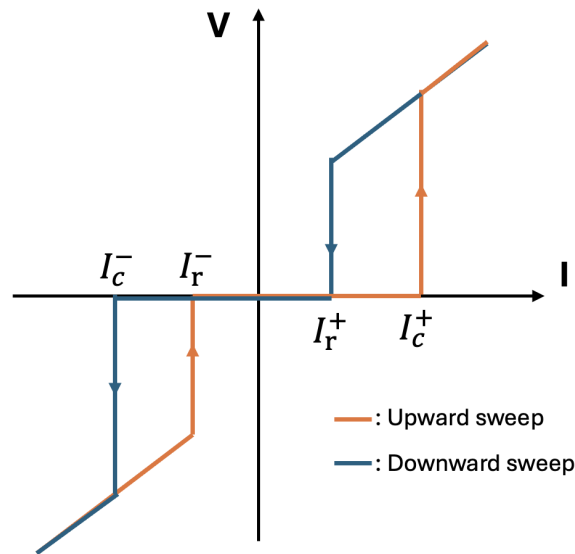


Figure II.1: Schematic showing definitions of the switching currents  $I_c^+$  and  $I_c^-$ . Also shown is the retrapping currents  $I_r^+$  and  $I_r^-$ .

Like the ordinary semiconductor diode, the superconducting diode is characterized by a nonreciprocity in charge transport. This means that the switching current will differ depending on the direction of current flow, leading to two different switching currents  $I_c^+$  and  $I_c^-$ . In particular, if the magnitude of the bias current is chosen to be between  $I_c^+$  and  $|I_c^-|$ , it enables supercurrent to flow in only one direction. As an example, assume  $I_c^+ > |I_c^-|$ , and pick a bias current  $|I_c^-| < I_{bias} < I_c^+$ . Applying this  $I_{bias}$  in the positive direction will lead to the current flowing as a supercurrent, as it is smaller than the positive switching current. Conversely, applying  $I_{bias}$  in the negative direction will lead to the material turning normal as the bias is larger than the negative switching current.

In general, nonlinear electrical responses such as the diode effect can occur when both time-reversal ( $\mathcal{T}$ ) and spatial inversion ( $\mathcal{P}$ ) symmetry is broken, an effect known as magnetochiral anisotropy (MCA) [29]. Inversion (time-reversal) symmetry means that the system retains the same properties under the transformation  $\mathbf{r} \mapsto -\mathbf{r}$  ( $t \mapsto -t$ ). In systems showing a finite SDE spatial inversion symmetry is often intrinsically broken, while time reversal symmetry often is broken by an applied magnetic field or through intrinsic magnetization [7]. An example can be seen in Ando's paper from 2020 [6]. There, the inversion symmetry was broken by stacking different types of materials, Nb,V,Ta, in an artificial superlattice without a center of inversion, and time reversal symmetry was broken by applying a magnetic field. Since that paper, it has been shown that the SDE is ubiquitous for thin film superconductors when applying an out of plane magnetic field [8]. This is due to inevitable differences at the edges of the sample, occurring from fabrication inaccuracies, which allow slightly different switching currents. Due to this effect, which is observable at out of plane fields as small as  $\mu_0 H \approx 300 \mu\text{T}$ , it is crucial to exclude it as a possibility when measuring diode efficiency at finite field. This is because the device might be misaligned by a few degrees in the refrigerator, leading to an out of plane field component when applying a field that is seemingly in-plane.

In order to quantify the SDE, it is common to introduce a the superconducting diode efficiency parameter  $\eta$ , defined as [8]:

$$\eta = \frac{I_c^+ - |I_c^-|}{I_c^+ + |I_c^-|}. \quad (\text{II.13})$$

This parameter takes values between -1 and 1, depending on the strength of the SDE, and will be used to quantify the SDE measurements in this thesis. A larger value of  $|\eta|$  means a larger diode effect, and a better superconducting diode. For the example current-voltage plot in Fig. II.1,  $\eta = 0$  since the positive and negative switching currents are equal. To estimate the error in measuring  $\eta$ , error propagation was used. Expanding  $\eta$  from Eq. (II.13), considering the nominator and denominator

as two separate variables, yields:

$$\eta \approx \eta^0 + \frac{\partial \eta}{\partial (I_c^+ + |I_c^-|)} (I_c^+ + |I_c^-|) + \frac{\partial \eta}{\partial (I_c^+ - |I_c^-|)} (I_c^+ - |I_c^-|). \quad (\text{II.14})$$

Taking the variance of both sides of Eq. (II.14) leads to, letting  $A := I_c^+ - |I_c^-|$ ,  $B := I_c^+ + |I_c^-|$ , and thus  $\eta = \frac{A}{B}$ :

$$\sigma_\eta^2 \approx \left| \frac{\partial \eta}{\partial A} \right|^2 \sigma_A^2 + \left| \frac{\partial \eta}{\partial B} \right|^2 \sigma_B^2, \quad (\text{II.15})$$

using the assumption that  $\text{Cov}(A, B) = 0$ . This is a valid assumption if the variance of the positive and negative switching currents are the same, which can be shown using the collinearity of the covariance:

$$\begin{aligned} \text{Cov}(A, B) &= \\ &= \text{Cov}(I_c^+ + |I_c^-|, I_c^+ - |I_c^-|) = \text{Cov}(I_c^+, I_c^+) - \text{Cov}(I_c^+, |I_c^-|) + \\ &\quad + \text{Cov}(|I_c^-|, I_c^+) - \text{Cov}(|I_c^-|, |I_c^-|) \\ &= \text{Var}(I_c^+) - \text{Var}(|I_c^-|) \end{aligned}$$

During measurements the variances of the positive and negative switching currents were comparable, indicating that this is a reasonable simplification. Continuing by inserting the partial derivatives into Eq. (II.15) yields:

$$\sigma_\eta^2 = \left( \frac{1}{B} \right)^2 \sigma_A^2 + \left( \frac{-A}{B^2} \right)^2 \sigma_B^2. \quad (\text{II.16})$$

Dividing both sides by  $\eta^2 = A^2/B^2$  and taking the square root:

$$\left( \frac{\sigma_\eta}{\eta} \right) = \sqrt{\frac{\sigma_A^2}{A^2} + \frac{\sigma_B^2}{B^2}} \quad (\text{II.17})$$

Multiplying both sides by  $\eta$ , and assuming  $\sigma_A = \sigma_B =: \sigma_I$ , the final error estimation is obtained:

$$\sigma_\eta = \eta \sqrt{\left( \frac{\sigma_I}{I_c^+ + |I_c^-|} \right)^2 + \left( \frac{\sigma_I}{I_c^+ - |I_c^-|} \right)^2} \quad (\text{II.18})$$

In the SDE calculations later in this thesis, the error bars show  $\pm 1\sigma_\eta$ , and the calculation of  $\sigma_I$  uses a 95 % confidence interval assuming a gaussian distribution of the switching currents.

## II.2 Magnetism

Magnetism in solid materials have traditionally been divided into two different types: ferromagnetism and antiferromagnetism [20]. Ferromagnets, which have been known for almost three thousand years, are characterized by their large net magnetization, where all magnetic moments lie in the same direction even in the absence of an applied field [30]. Antiferromagnets on the other hand have been known only for about a century, owing to their vanishing net magnetization that makes them behave similarly to non-magnetic materials. Recently, however, there have been theoretical developments delimiting a third distinct type of magnetism, altermagnetism [20].

This section begins with a short overview of ferromagnetism in order to provide a basis for comparison with the magnetic materials used in this thesis. How ferromagnetism interacts with superconductors will also be briefly covered to illustrate how a magnetic material might influence a nearby superconductor. Afterwards an overview of antiferromagnetism and altermagnetism will be given, along with a discussion on how these phenomena differ from and are alike to ferromagnetism.

### II.2.1 Ferromagnetism

According to the Bohr-van Leeuwen theorem, classical physics alone cannot explain the net magnetization found in ferromagnetic materials [30]. To illustrate this, one can consider the magnetic dipole-dipole interaction, which might be expected to be the fundamental mechanism underlying the magnetization of materials if only classical physics were considered. Two magnetic dipoles  $\boldsymbol{\mu}_1, \boldsymbol{\mu}_2$  separated by a distance  $r$  has an energy of [30]:

$$E_{\text{dip}} = \frac{\mu_0}{4\pi r^3} \left( \boldsymbol{\mu}_1 \cdot \boldsymbol{\mu}_2 - \frac{3}{r^2} (\boldsymbol{\mu}_1 \cdot \mathbf{r})(\boldsymbol{\mu}_2 \cdot \mathbf{r}) \right). \quad (\text{II.19})$$

The order of magnitude of this energy for two electrons in a solid can be estimated by letting  $\mu_1 = \mu_2 \approx \mu_B$  and  $r \approx 1 \text{ \AA}$ , which gives  $E_{\text{dip}}/k_B \approx 1 \text{ K}$ . This can be compared to the temperature at which magnetization breaks, which can be on the order of  $\sim 1000 \text{ K}$ . An example magnetization graph for a material with a much larger critical temperature is shown in Fig. II.2, where the magnetization is suppressed at a temperature of  $T_c = 768 \text{ K}$ . This means that in a classical picture, the energy gained in aligning magnetic moments is much smaller than the temperatures at which this alignment breaks.

Thus, one needs quantum physics to describe magnetism, and the ingredients needed are Pauli's exclusion principle as well as Coulomb repulsion between electrons [32]. Consider the case of two electrons in a box, in two different states  $a$  and  $b$ . Due to the Pauli exclusion principle the total wavefunctions for the electrons must be antisymmetric. This leads to two options: either an antisymmetric singlet state

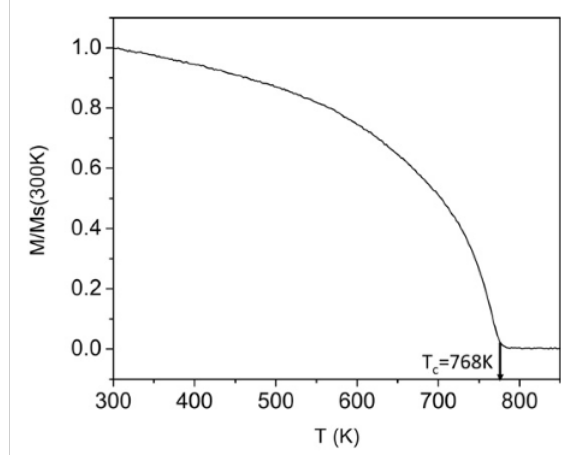


Figure II.2: Magnetization (normalized with the room temperature magnetization) as a function of temperature for a  $\text{Ni}_{83}\text{Fe}_{13}$  thin film. Magnetization breaks at  $T_c = 768\text{ K}$ , much larger temperatures than  $E_{\text{dip}}/k_B \approx 1\text{ K}$ . This illustrates that there must be some mechanism other than magnetic dipole interaction causing the magnetization. Figure taken from [31].

$\chi_S$  together with a symmetric spatial state or a symmetric triplet state  $\chi_T$  together with an antisymmetric spatial state. Thus, with the first electron at position  $\mathbf{r}_1$  and the second at  $\mathbf{r}_2$ , the wavefunctions for the singlet and triplet state can be written, respectively:

$$\begin{cases} \Psi_S = \frac{1}{\sqrt{2}} [\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) + \psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1)] \chi_S \\ \chi_S = \frac{1}{\sqrt{2}} (|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle) \\ \Psi_T = \frac{1}{\sqrt{2}} [\psi_a(\mathbf{r}_1)\psi_b(\mathbf{r}_2) - \psi_a(\mathbf{r}_2)\psi_b(\mathbf{r}_1)] \chi_T \\ \chi_T = |\uparrow\uparrow\rangle, |\downarrow\downarrow\rangle, \frac{1}{\sqrt{2}} (|\uparrow\downarrow\rangle + |\downarrow\uparrow\rangle) \end{cases} \quad (\text{II.20})$$

with the energies  $E_S = \langle \Psi_S | \hat{H} | \Psi_S \rangle$  and  $E_T = \langle \Psi_T | \hat{H} | \Psi_T \rangle$ . For two electrons, with spins  $\hat{\mathbf{S}}_1$  and  $\hat{\mathbf{S}}_2$  and  $\hat{\mathbf{S}}_{\text{tot}} := \hat{\mathbf{S}}_1 + \hat{\mathbf{S}}_2$ , it follows that  $(\hat{\mathbf{S}}_{\text{tot}})^2 = \hat{\mathbf{S}}_1^2 + \hat{\mathbf{S}}_2^2 + 2\hat{\mathbf{S}}_1 \cdot \hat{\mathbf{S}}_2$ . From ordinary quantum mechanics, combining two spin  $\frac{1}{2}$  particles results in their sum having spin quantum number  $s$  equal to 0 or 1. Since an angular momentum operator  $\hat{\mathbf{L}}^2$  has an eigenvalue of  $l(l+1)$  it follows that  $\hat{\mathbf{S}}_1 \cdot \hat{\mathbf{S}}_2$  has the eigenvalues  $\frac{1}{2} \left( s(s+1) - 2 \cdot \frac{3}{4} \right) = -\frac{3}{4}$  or  $\frac{1}{4}$  in the case of the singlet ( $s=0$ ) and triplet ( $s=1$ ) state respectively. Using this, one can write the Hamiltonian in the form of an effective Hamiltonian:

$$\hat{H} = \frac{1}{4} (E_S + 3E_T) - 2J \hat{\mathbf{S}}_1 \cdot \hat{\mathbf{S}}_2, \quad (\text{II.21})$$

where the exchange integral  $J$  is defined as  $J := \frac{1}{2}(E_S - E_T)$  [30]. When this exchange integral is larger than zero, the triplet state is favored, and when it is smaller than zero, the singlet state is favored. In the considered case of two electrons in a box, the triplet state will minimize the energy cost from Coulomb repulsion, and thus be

the favored state [32].

Generalizing this model to the case of many electrons can be done by letting  $\hat{H} = -\sum J_{ij} \hat{\mathbf{S}}_i \cdot \hat{\mathbf{S}}_j$ , where  $J_{ij}$  is the exchange integral between sites  $i$  and  $j$ , often taken to be constant for nearest-neighbour interactions and zero otherwise [30]. For ferromagnets the exchange constant  $J$  is positive, favoring the triplet state and thus the parallel spin states.

How a ferromagnet interacts with a nearby superconductor has been studied for many years. Originally, the interest lay in the suppression of superconductivity due to singlet-state repulsion. An ordinary superconductor has, as described in Section II.1, electrons of opposite spin and momentum pair up and form Cooper pairs. In the setting of a superconductor close to a ferromagnet, assuming a strong enough exchange interaction, a given spin up electron A will be mostly surrounded by other spin up electrons. Since the attractive phonon force is very short-ranged, a spin down electron B attempting to lower its energy by coupling to A and forming a Cooper pair first has to travel through a large region of unfavorably oriented spins. This effective shielding of electrons with opposite spin suppresses the ordinary phonon-mediated superconductivity [33]. Though ferromagnetism suppresses this ordinary superconductivity, there are cases where ferromagnets at the interface of superconductors may lead to exotic types of superconductivity with spin-triplet Cooper pairs [13].

## II.2.2 Antiferromagnetism and CuMnAs

If a ferromagnet arises from a positive exchange integral  $J > 0$  (Eq. (II.21)), then an antiferromagnet arises from a negative exchange integral  $J < 0$  [30]. In this case, it is favorable for nearest neighbours to have opposite magnetic moments, which results in net zero magnetization for the material as a whole. An antiferromagnet can be viewed as two interpenetrating sublattices, each with the same magnetic moment on all sites, as illustrated in Fig. II.3. Since the total magnetization  $\mathbf{M}$  is zero, it cannot be used as an order parameter as it is for ferromagnets. Instead, one uses the difference between the magnetizations of the two different sublattices  $\mathbf{M}_1$  and  $\mathbf{M}_2$ , the Néel vector, as the order parameter:  $\mathbf{L} \propto \mathbf{M}_1 - \mathbf{M}_2$  [30, 34]. This ordering holds below a critical temperature, called the Néel temperature  $T_N$  [30]. In general, it is hard to flip the Néel vector using magnetic fields, due to the net zero magnetization of AFMs.

CuMnAs is a semimetal [35], and a so-called high temperature antiferromagnet, with a Néel temperature of  $T_N = 480$  K [15]. The magnetic moments in the material are supplied by the Mn atoms [18], see Fig. II.4 for an illustration of the alignment of the magnetic moments in tetragonal CuMnAs. This figure shows how the Mn atoms' magnetic moments compensate each other, making the net magnetization of the unit cell equal to zero, compare also to Fig. II.3. For thin film CuMnAs,

the Néel vector has been shown to flip at applied magnetic fields on the order of 2 T [36]. It has also been demonstrated that the Néel vector of CuMnAs can be controlled and read out using electrical pulses, taking advantage of relativistic spin-orbit coupling effects [18]. The applied current pulses generate magnetic fields of opposite polarity inside the unit cell, flipping the Néel vector. In relation to the earlier discussion about broken time- and inversion symmetries in Section II.1.2, CuMnAs breaks both time- and inversion symmetry [37]. However, the combined time and inversion symmetry  $\mathcal{PT}$  is preserved.

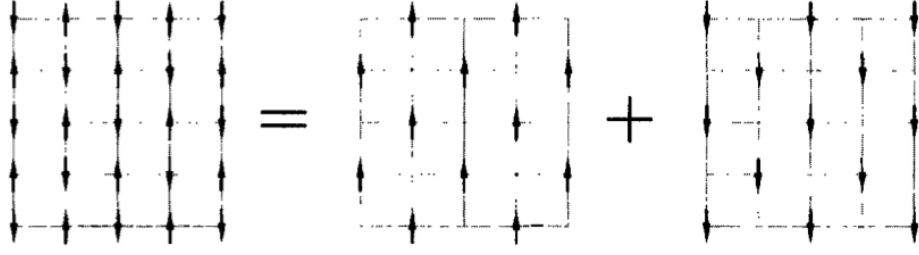


Figure II.3: Decomposition of an antiferromagnetic lattice into two sublattices where all sites on the same sublattice have equal magnetic moments. Figure taken from [30].

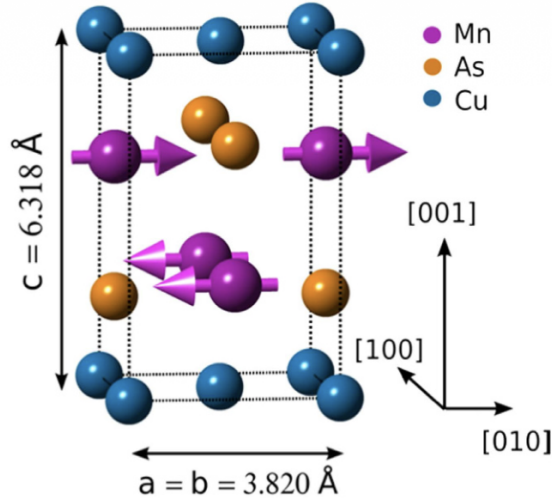


Figure II.4: Unit cell illustrating the structure of tetragonal CuMnAs, where the antiferromagnetic alignment of the magnetic moments is visible. Also shown is that the magnetic moments are due to the Mn atoms. Later in the thesis the [001]-direction will coincide with the out of plane direction  $x$ . Figure taken from [19].

### II.2.3 Altermagnetism and MnTe

Altermagnetism is a recently discovered type of magnetism [38], with similarities both to ferromagnetism and antiferromagnetism. Like antiferromagnetism, the

macroscopic net magnetization of an altermagnet is zero, and the order parameter is the Néel vector [21]. What differs from an antiferromagnet is that the opposite-spin sublattices are connected, not via translation or inversion symmetry as in an antiferromagnet, but with specific crystal rotations [20]. This leads to spin-split bands in momentum space, which is also characteristic of ferromagnets. The ferromagnetic spin-splitting, however, comes from the net magnetization in the material breaking time-reversal symmetry [20]. The spin-splitting for altermagnets also differs from ferromagnetic spin-splitting in that  $E_{\uparrow}(\mathbf{k}) - E_{\downarrow}(\mathbf{k}) = E_{\downarrow}(-\mathbf{k}) - E_{\uparrow}(-\mathbf{k})$ , that is, the spin splitting alternates sign. For ferromagnets one spin band is shifted by a constant value from the other. See Fig. II.5 for an illustration of this and a comparison between the altermagnetic, antiferromagnetic, and ferromagnetic cases.

The material used in this thesis, MnTe, has been experimentally confirmed as having the spin-split structure characteristic of an altermagnet by photoemission spectroscopy [14]. It is as of yet unclear how stable this altermagnetic phase is, but on the basis of this experimental evidence it will be referred to as an altermagnet in this thesis. It is a semiconductor with a band gap of  $\approx 1.4$  eV [21].

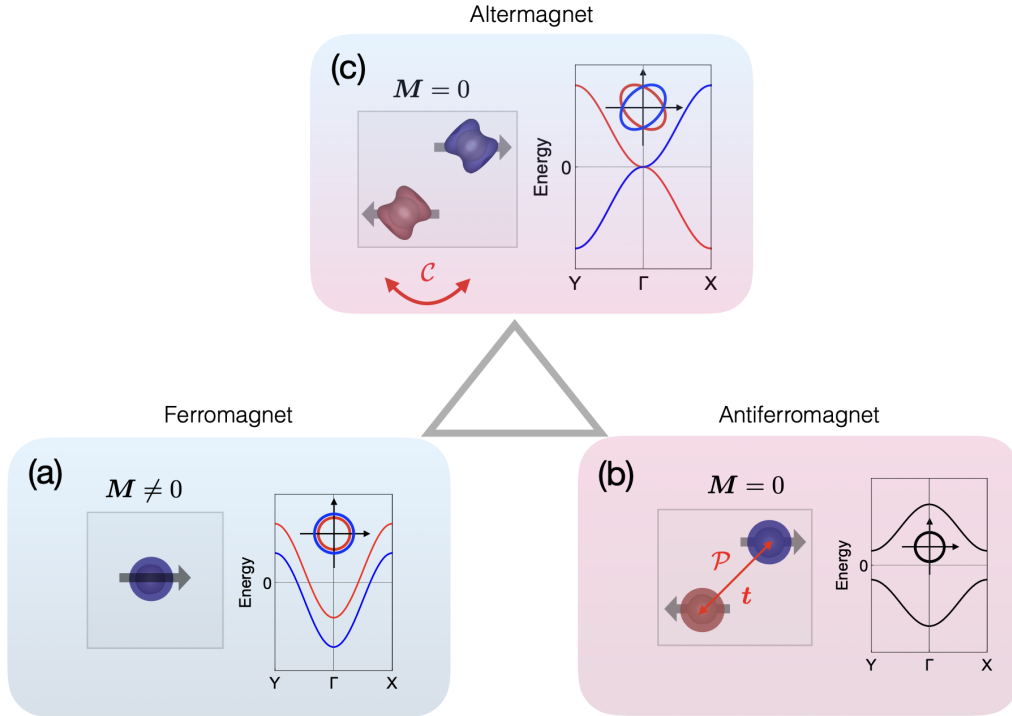


Figure II.5: Illustration of the differences between ferromagnets, antiferromagnets, and altermagnets. (a) Ferromagnet, with a net magnetization and spin-split band structure. One spin band is shifted by a constant amount relative to the other. (b) Antiferromagnet, net zero magnetization and opposite-spin sublattices connected by inversion ( $\mathcal{P}$ ) or translation ( $\mathbf{t}$ ). Spin-degenerate bands. (c) Altermagnet, with opposite-spin sublattices connected by crystal rotations ( $\mathcal{C}$ ). Alternating sign of the spin-splitting. Image taken from [20].

## II.2.4 Kondo Effect

When cooling down non-superconducting metals, the resistance decreases as the temperature reduction reduces scattering events. However, the resistance saturates when the temperature drops below about 10 K, due to static defects [39]. This behavior changes dramatically, however, when magnetic atoms are present in the material in addition to the metal. In this case, the resistance continues to increase as the temperature is lowered, see Fig. II.6. Also shown in the figure is the Kondo temperature,  $T_K$ , which is the temperature at which the resistance is at a minimum. A relevant example for this thesis is the alloy Mn in Al, which has a  $T_K$  of 500 K [40].

This effect originates from interactions between the magnetic moments of the ions and the magnetic moments of the conduction electrons. The electrons form a cloud around the magnetic impurity, with opposite spin to it. This forms a quasi-bound state [30]. This screening of the magnetic moment of the impurity leads to an increase in scattering which increases resistance. At low temperatures the resistance is proportional to  $-\log T$  [30].

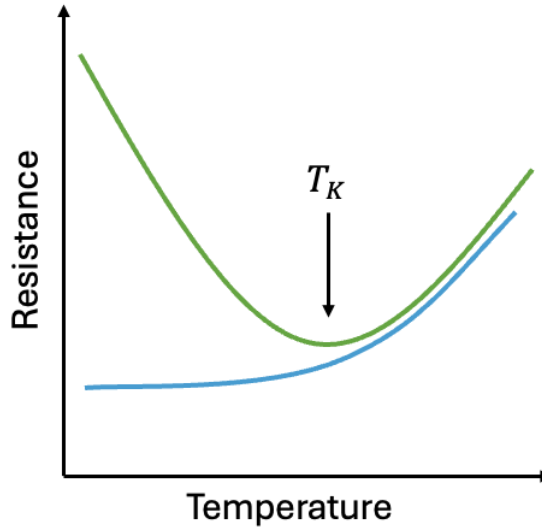


Figure II.6: Green curve shows the Kondo effect, leading to the resistance increasing as the temperature is lowered. Blue curve shows the resistance of a non superconducting metal saturating when the temperature is lowered. The Kondo temperature  $T_K$  is also shown at the minimum of the green curve.

# III

## MnTe/Al

---

This chapter investigates bar devices of the altermagnetic-superconductor (MnTe-Al) heterostructure. Material growth was made by Peter Wadley's group at the University of Nottingham. Fabrication of the devices was done by Lucas Casparis (University of Copenhagen). Measurements were performed by me.

A scanning electron microscope (SEM) image of a cross-section of the wafer can be seen in Fig. III.1. From this image the thicknesses of the Al and MnTe layers are estimated to be 12 nm and 60 nm, respectively.

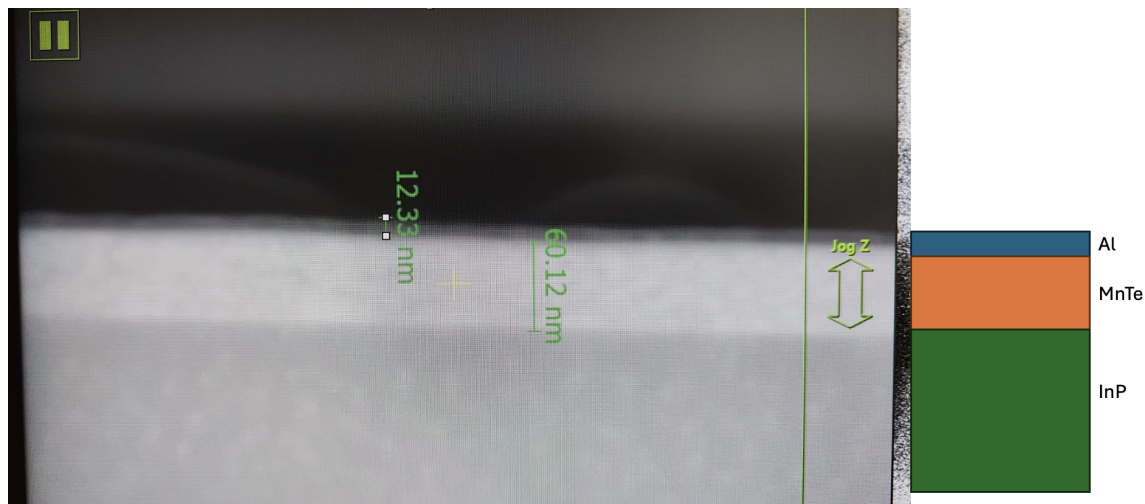


Figure III.1: SEM of the MnTe/Al wafer used in this thesis, as well as a schematic of the materials. Image taken by Filip Krizek (Czech Academy of Sciences).

### III.1 Measurement Setup

The device measured was designed to have several bar devices at different angles to each other. An image of the wire bonded device as well as a schematic showing

which number corresponds to which pad can be seen in Fig. III.2. The length  $l$  of the channels were varying, and the width  $w$  of each channel was  $1\text{ }\mu\text{m}$ . A summary of the aspect ratios of the channels between different pads, defined as  $l/w$ , can be seen in Table III.1.

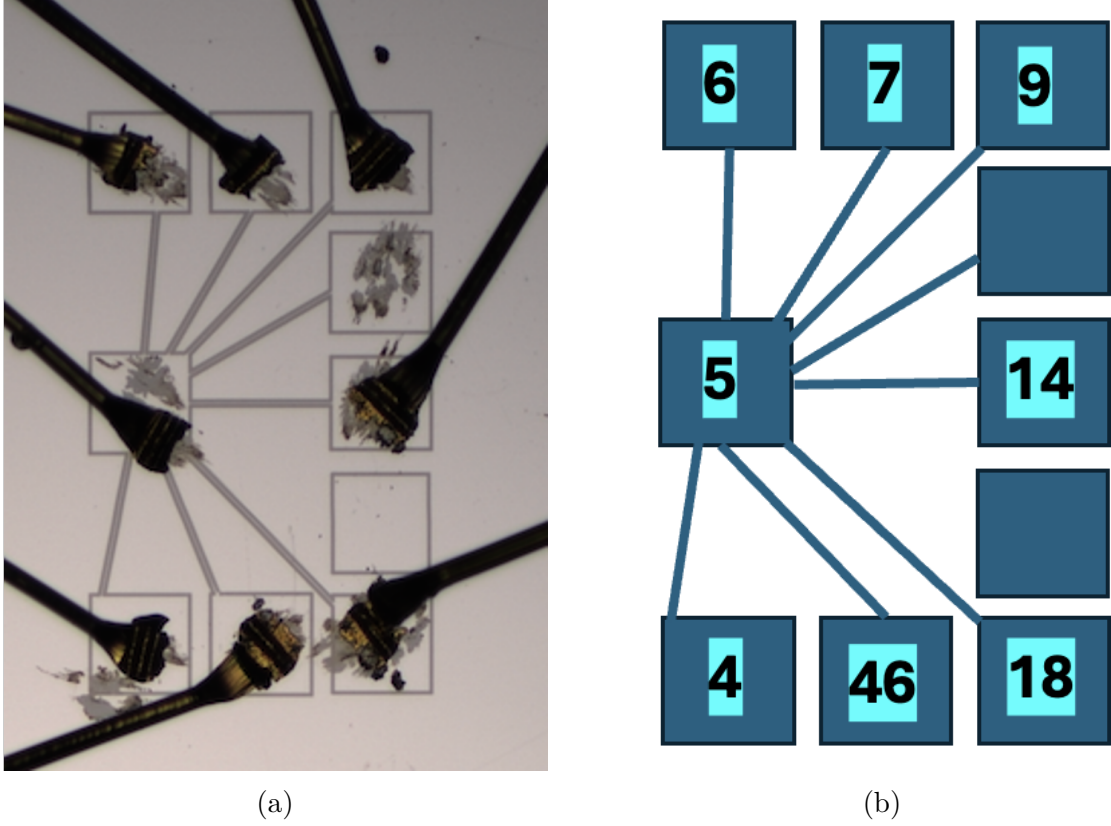


Figure III.2: (a) Optical microscopy image of the bonded device. (b) Schematic showing the pin number of the cryostat connected to the corresponding pad.

| Pad x – Pad y | Aspect Ratio |
|---------------|--------------|
| 5–6           | 59.5         |
| 5–7           | 68.6         |
| 5–9           | 88.2         |
| 5–14          | 59.3         |
| 5–18          | 86.8         |
| 5–46          | 64.0         |
| 5–4           | 60.9         |

Table III.1: Aspect ratios for the different channels in the device, defined as  $l/w$ , where  $l$  is the length of the channel and  $w$  is the width. Since all devices had  $w = 1\text{ }\mu\text{m}$ , the aspect ratios are also the channels' lengths in  $\mu\text{m}$ . See Fig. III.2b for the placement of each pad on the device.

The measurements were done in a Kiutra L-rapid cryostat, which can sustain a temperature of 300 mK indefinitely and a temperature of 100 mK for up to three

hours. Measurements were performed in a two-terminal setup using lock-in amplifiers and a Basel Precision Instrument current-to-voltage converter. When measuring between two pads, current was biased through one of the pads, and the other pad was grounded. The other pads were floated, so that current would flow only between the biased and the grounded pads. See Fig. III.3 for a schematic showing the case of measuring between pads 5 and 14. When applying bias in this way, the Basel IV-converter divides the voltage from the lock-in output by 100 and applies it to the other channel on the same side. That is, if the lock-in output is set to 100 mV in Fig. III.3, there will be a voltage difference of 1 mV between pads 5 and 14. By controlling the voltage in this way and measuring the current via the lock-in input, the differential resistance could be calculated. This method was chosen because when testing the connections it was found that there was a path to ground from all pads. Based on Fig. III.2a, the path to ground was determined to most likely be from the bonding at pad 5. This is because the gold bonding wire can be seen to touch both pad 5 and the rest of the chip, creating a path for current between pad 5 and the rest of the chip. As the other pads were connected to pad 5, this meant that all pads were connected to ground.

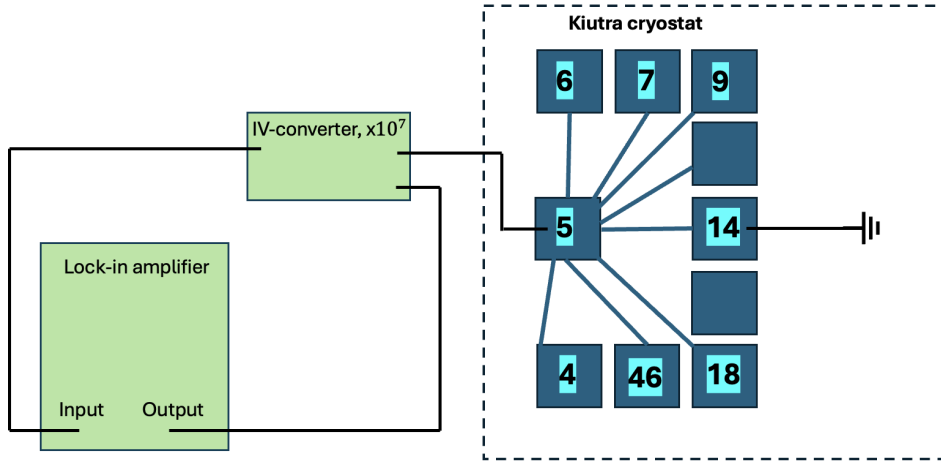


Figure III.3: Schematic showing the measurement setup in the case of measuring between pads 5 and 14 on the device. Bias voltage is applied by the lock-in amplifier via the IV-converter, and the other pad is grounded. The current running through the device is amplified by  $10^7$ , converted into a voltage, and read out by the same lock-in providing the bias.

## III.2 Measurement Results

In order to compare the different channels the sheet resistance was used, due to the different aspect ratios. In the Kiutra the line resistances were known to be  $3.4 \text{ k}\Omega$ . The sheet resistance was thus calculated as  $R_S = (R_{measured} - 3.4 \text{ k}\Omega) / (\text{Aspect ratio})$ ,

using the aspect ratios from Table III.1.

To check whether the device would turn superconducting, a constant voltage bias of  $50\text{ }\mu\text{V}$  was applied over the channel which was to be measured. The current through the sample was measured and from these the sheet resistance as a function of temperature in the range  $0.3\text{ K}$  to  $3.5\text{ K}$  was measured. The results from these measurements can be seen in Fig. III.4, where it is clear that the sample never turned superconducting in the measured range. In fact, the resistance increased as the temperature was lowered, characteristic of an insulator or semiconductor. A possible explanation for the absence of superconductivity is that the magnetic moments in MnTe suppressed the superconductivity. The much larger sheet resistance seen between pads 5 and 46 is not believed to be physical, and might be explained by variations in the fabrication. It can also be noted that the sheet resistance varies up to a factor of 6 between the different pads, even when excluding the one between pads 5 and 46, which might indicate directional dependence on resistivity of the heterostructure. It could however also be due to differences in fabrication.

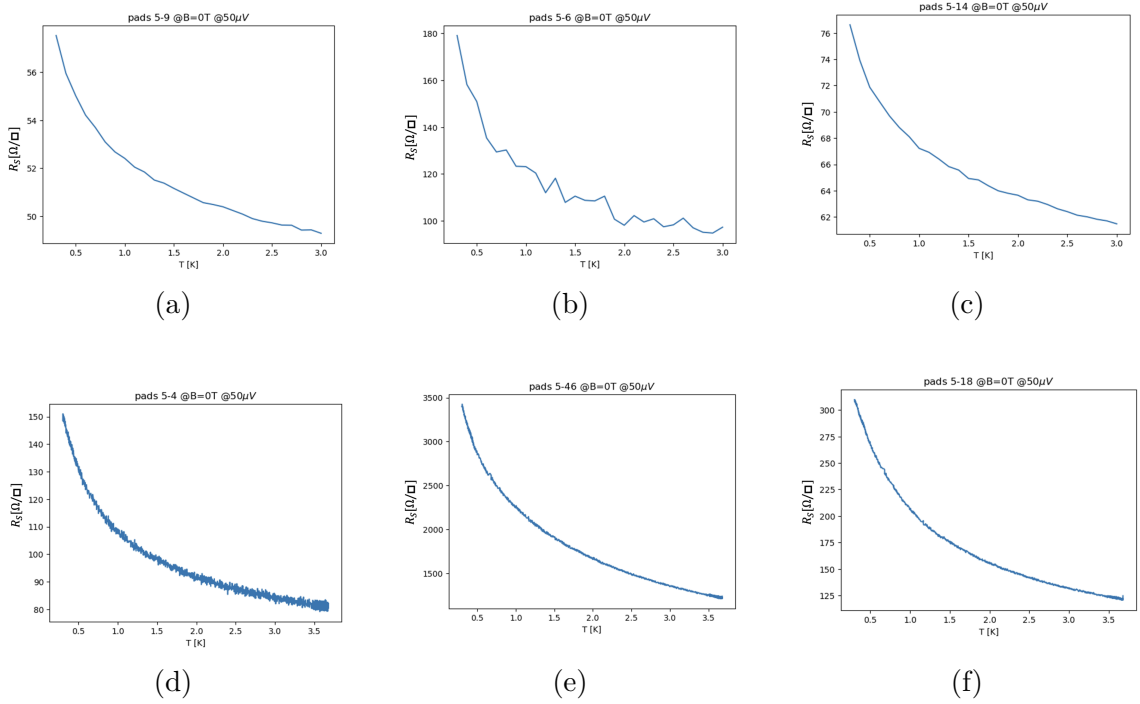


Figure III.4: Sheet resistance as a function of temperature for six of the channels, at a constant  $50\text{ }\mu\text{V}$  bias in the temperature range  $0.3\text{ K}$  to  $3.5\text{ K}$ . Measured between pads (a) 5 and 9, (b) 5 and 6 (c) 5 and 14, (d) 5 and 4, (e) 5 and 46, (f) 5 and 18. (a)-(c) are temperature up sweeps, while (d)-(f) are temperature down sweeps.

One possible explanation for the plots in Fig. III.4 is that MnTe dominated the charge transport. If this were the case we would expect the conductance to depend mainly on the charge carrier density since MnTe is a semiconductor. In a semiconductor the charge carrier density is generally proportional to  $\exp \frac{-E_g}{2k_B T}$  [41], which leads to  $\log R \propto 1/T$ . To check whether this would be consistent with

our measurements,  $\log R$  was plotted against  $1/T$ , see Fig. III.5. Since this plot is not linear, the resistance increase with lower temperatures cannot be explained by semiconductor-like charge carrier depletion and must be explained in some other way.

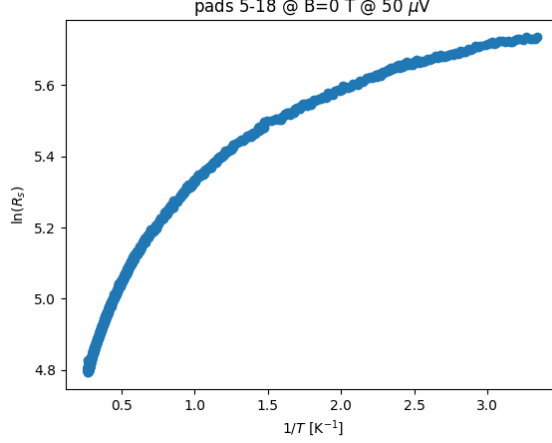


Figure III.5:  $\log R$  as a function of  $1/T$ , from the data in Fig. III.4f.

Possible explanations to the  $R(T)$  behavior seen in Fig. III.4 could be the Kondo effect discussed in Section II.2.4, see also Fig. II.6. A possibility is that some of the Mn atoms inadvertently formed an alloy with the aluminum during growth of the material. Since the Kondo temperature is around 500 K for Mn in Al samples, this could be consistent with our findings. In order to check whether the data fits the  $R \propto -\log T$  expected from the Kondo effect,  $R$  was plotted against  $\log T$  and a linear fit was performed, see Fig. III.6. From this figure it can not be ruled out that  $R \propto -\log T$ . The analysis for the data between the other pads looked similar as well.

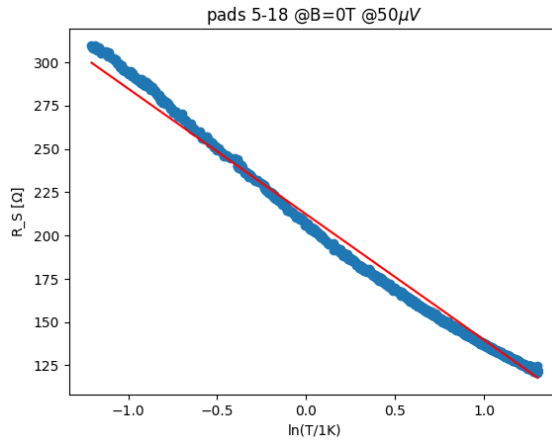


Figure III.6:  $R$  as a function of  $\log T$ , using the data from Fig. III.4f. The red line is a linear fit over the entire range.

Another possible explanation for the behavior seen is weak localization. This is a quantum effect in which phase-coherent multiple scattering leads to an increased

probability of backscattering for electrons and thus increased resistance. This effect is temperature dependent because lowering the temperature increases the phase coherence length [41]. Like the Kondo effect, weak localization can also lead to a resistance which increases logarithmically with decreasing temperature. This would give a linear dependence in an  $R$  vs.  $\log T$  plot, and can thus also not be ruled out on the basis of Fig. III.6.

The next steps to investigate this would be to reproduce the results with directional dependent resistance and increasing resistance at low temperatures. A structural analysis could also be performed to investigate the possibility of alloying, in order to determine whether the  $R(T)$  dependence is due to the Kondo effect. Since there was no superconductivity in the chip, it was decided to investigate the CuMnAs/Al heterostructure instead.

# IV

## CuMnAs/Al

---

This chapter investigates bar devices of the antiferromagnetic-superconductor (CuMnAs/Al) heterostructure. Growth was done by Filip Krizek (Czech Academy of Sciences) and fabrication was done by Lucas Casparis (University of Copenhagen). Measurements and data analysis were done in collaboration with August Jacobi, see his thesis for more on this material stack.

### IV.1 Growth and Fabrication

The antiferromagnetic CuMnAs used in this thesis was grown by Filip Krizek as detailed in the 2020 paper [16], with an epitaxial layer of CuMnAs on top of a GaP substrate. In addition to the described procedure in that article, the material studied in this thesis was capped with a thicker layer of aluminum. A transmission electron microscopy (TEM) image of the material used in this thesis can be seen in Fig. IV.1, where it can additionally be seen that a layer of AlAs has formed at the intersection between the Al and CuMnAs. To confirm that the material was antiferromagnetic, a magnetization measurement was done with a SQUID by Filip Krizek, taken at different points on the wafer. The result of this measurements can be seen in Fig. IV.2, and should be compared to the free Mn magnetization of  $\sim 5 \mu_B$ , where  $\mu_B = \frac{e\hbar}{2m_e}$  is the Bohr magneton [42]. Since the magnetic moments in CuMnAs comes from the Mn atom, this measurement indicates that the magnetic moments from Mn are strongly suppressed or compensated in this growth. This measurement was performed on an identically grown wafer to the one used in this thesis. Fabrication and bonding of the devices was done by Lucas Casparis, and the chip included several bar devices at different angles. In this thesis, the focus was on studying two of the bar devices, which were oriented perpendicularly to one another. Optical microscopy of the devices can be seen in Fig. IV.3 and Fig. IV.4 for devices C2 and F1 respectively.

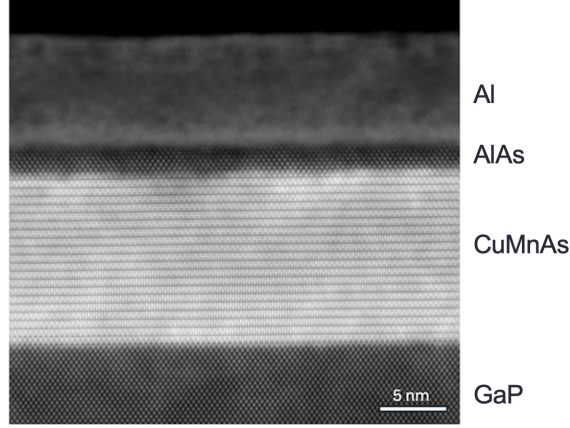


Figure IV.1: TEM cross-section image of the CuMnAs/Al wafer used in this thesis. Image taken by Filip Krizek at the Czech Academy of Sciences, Prague.

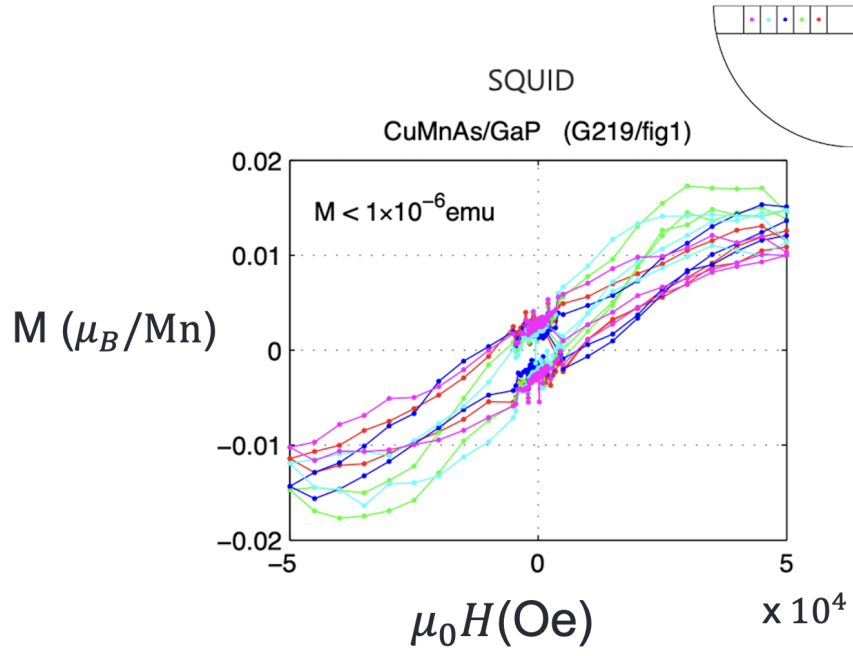


Figure IV.2: Magnetization measurement with a SQUID as a function of applied field for different positions on an identically grown wafer. y-axis is in units of Bohr magneton per manganese atom. Measurements taken by Filip Krizek.

## IV.2 Measurement Setup

The chip was measured using a cryofree dilution refrigerator (Oxford Instruments, Triton) with a base temperature of 20 mK. All measurements, with the exception of the  $T_c$  measurement, were carried out at base temperature. The measurements were done using dc current, with current bias generated by applying a voltage from a commercially available Source Measure Unit (Keithley 2614B) over a resistance of 200 k $\Omega$  for one of the devices and 199.6 k $\Omega$  for the other device. The reason for this was that only one 200 k $\Omega$ -resistor was found. The voltage and current were

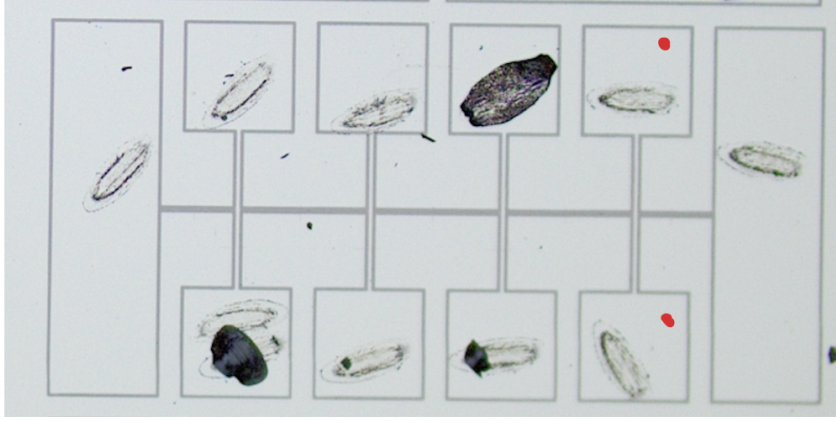


Figure IV.3: Optical microscopy image of device C2. Measurements done between the bond pads marked with red dots. Note that the rest of the device is not connected to the part measured in this thesis. Image taken by Lucas Casparis.

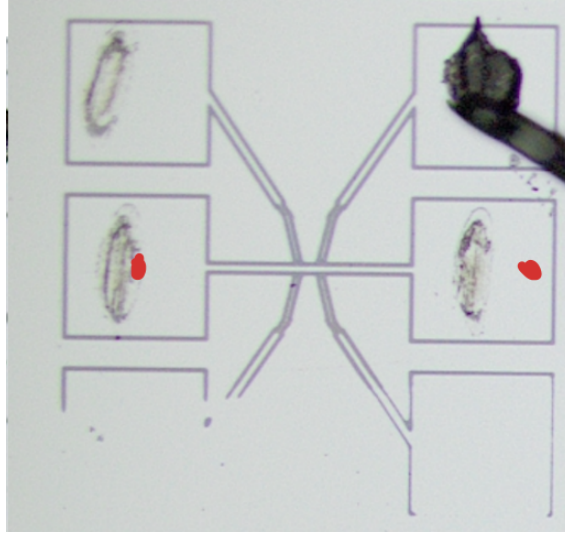


Figure IV.4: Optical microscopy image of device F1. Measurements done between the bond pads marked with red dots. Note that the rest of the device is not connected to the part measured in this thesis. Image taken by Lucas Casparis.

measured using Digital Multimeters (Agilent 34401A). The current was amplified using a current-to-voltage converter from Basel Precision Instruments, with a gain of  $10^5$  and a cutoff frequency of 3 kHz. The voltage measured was amplified using a preamplifier from Basel Precision Instrument with a gain of  $10^3$ . A schematic of the measurement setup on one of the bar devices can be seen in Fig. IV.5. The setup was identical for the two devices except for the bias resistance. The digital multimeters were read out using a PC. Magnetic fields were applied using a three-axis vector magnet, capable of applying a magnetic field of up to 6 T in the  $z$ -direction, and up to 1 T in the  $x$ - and  $y$ -directions. Both devices lay in the  $yz$ -plane, and had a length of  $180\text{ }\mu\text{m}$  and width of  $5\text{ }\mu\text{m}$ . Current flow in C2 (F1) was parallel to the  $z$ -axis ( $y$ -axis), see Fig. A.1. The crystallographic direction  $[001]$  is the out of plane

direction, parallel to  $H_x$  in Fig. A.1, see also Fig. II.4.

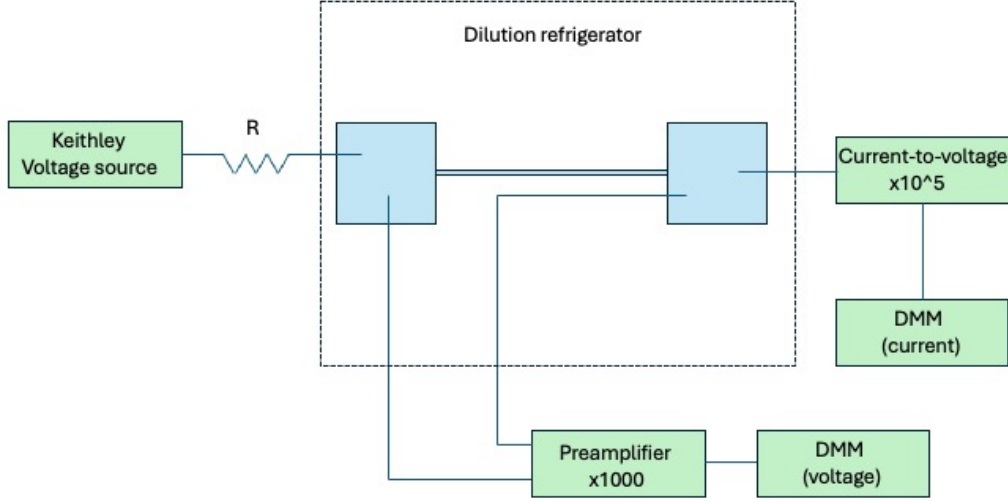


Figure IV.5: Schematic of the measurement setup showing how the bar devices were measured. A current bias was applied by applying a voltage over the bias resistor  $R$ . For C2,  $R = 199.6 \text{ k}\Omega$  while for F1,  $R = 200 \text{ k}\Omega$ .

### IV.3 Measurement Results

#### IV.3.1 Basic Characterization

Sample I-V curves for C2 and F1 can be seen in Fig. IV.6. In these figures, the flat voltage region indicates superconductivity, see also Fig. II.1. The devices have similar differential resistances of around  $200 \Omega$  in the normal state, independent of applied magnetic field. The devices show a difference in the size of their switching currents, in the case of Fig. IV.6 by around a factor of 5. The switching current of F1 was consistently smaller than the switching current of C2. The reason for this could be due to variations in fabrication, or differences in the crystallographic direction of the CuMnAs. Also note that the retrapping current is much smaller than the switching current. This can be seen by noting that the transition from normal to superconducting happens at negative currents quite close to zero, while the transition to the normal state from the superconducting state happens at around  $16 \mu\text{A}$  for C2 and at around  $3 \mu\text{A}$  for F1. This is attributed to heating, as discussed in Section II.1.2. The critical temperature of the devices were measured to be  $\approx 250 \text{ mK}$  for C2, and  $\approx 175 \text{ mK}$  for F1, see Fig. IV.7. The resistance in this figure was estimated by a linear fit between  $0 \mu\text{A}$  to  $1 \mu\text{A}$  for C2 and between  $0 \mu\text{A}$  to  $0.5 \mu\text{A}$  for F1. This should be compared with  $T_c$  for bulk Al at  $1.2 \text{ K}$ , and with Al thin films

where  $T_c$  can reach values of up to 2.4 K [43]. The CuMnAs thus seems to degrade the superconductivity in Al, which indicates some kind of interaction between the materials.

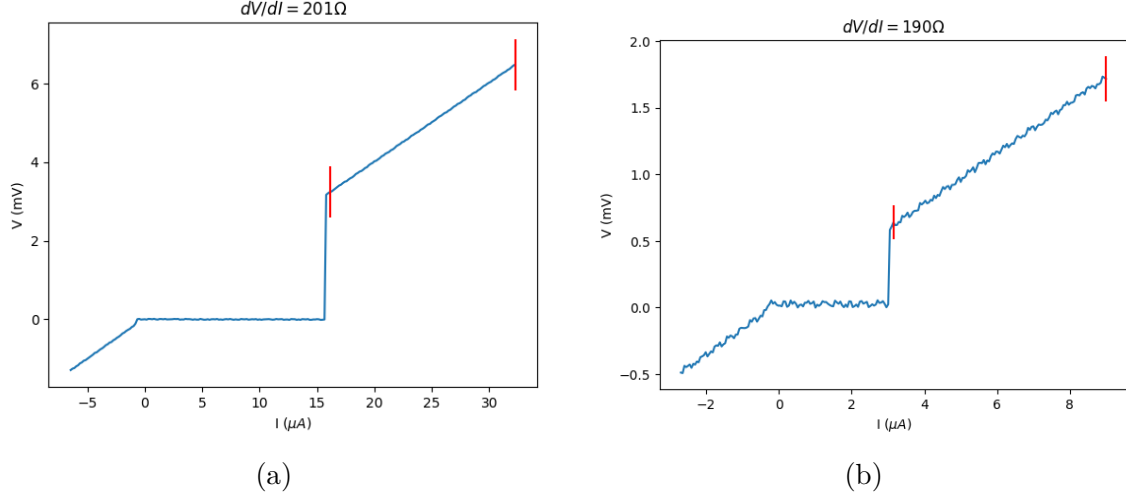


Figure IV.6: Sample I-V curves for C2 and F1, measured at a temperature of approximately 20 mK with zero applied magnetic field. The current was swept from the negative to positive side. A linear fit in the normal state region marked with vertical red bars yielded a differential resistance of  $dV/dI = 201 \Omega$  for C2 and  $dV/dI = 190 \Omega$  for F1.  $I_c \approx 16 \mu A$  for C2 and  $I_c \approx 3 \mu A$  for F1.

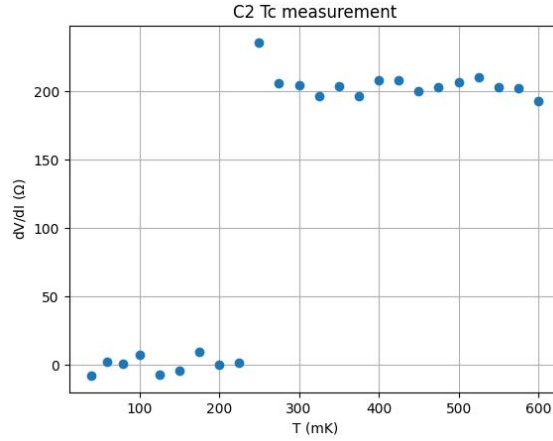


Figure IV.7: Differential resistance as a function of temperature for C2 in the range 40 mK to 600 mK. Differential resistance calculated via linear fit of  $V(I)$  in the range 0  $\mu A$  to 1  $\mu A$  for C2.  $T_c \approx 250$  mK for C2.

To determine the critical out of plane field for the devices,  $I_c(H_x)$  was measured. The results can be seen in Fig. IV.8 for the two field sweep directions, for the positive switching currents. A small offset can be seen for both devices, where the peak of the up sweep is at a slightly more positive field than for the down sweep. It can also be noted that the critical field is larger for C2 than F1, around 7.5 mT compared to  $\approx 3$  mT.

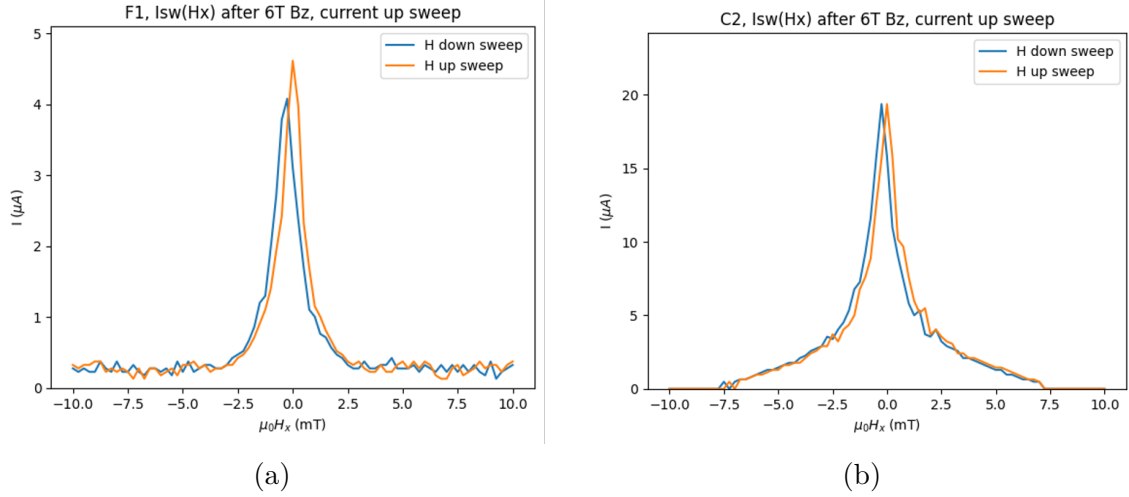


Figure IV.8:  $I_c(H_x)$  for F1 and C2. Shown in the figure are the results from the down (blue) and up (orange) field sweeps. A small offset is visible in the peak for both devices. That the switching current for F1 never reaches zero is an artifact, it goes to zero at  $\mu_0 H_x \approx \pm 3$  mT. Resolution in  $\mu_0 H_x$  was 250  $\mu$ T.

During the measurements it was found that there were two parameters that strongly affected the measurements: delay time, defined as the time spent measuring each current step, and number of sweeps. The effect that the number of sweeps has on the devices is illustrated in Fig. IV.9, where 50 I-V sweeps from  $-25 \mu\text{A}$  to  $25 \mu\text{A}$  was taken for C2 directly after a field excursion to  $\mu_0 H_y = 1$  T. In this example one can see the switching current relax after  $\sim 20$  sweeps. F1 showed similar behavior. The more important parameter, however, was the delay time. This was shown after a delay time test was done, performing 10 IV-sweeps with 0.01, 0.02, 0.05, 0.1, 0.2, and 0.5 s delay times, see an example in Fig. IV.10. The average switching current from the 10 sweeps was extracted for each delay time, and the results can be seen in Fig. IV.11. From this figure it can be seen that F1 shows a stronger dependence than C2 on the delay time, with superconductivity almost disappearing completely when the delay time goes toward zero. From these tests it was decided to use a 0.2 s delay time, as a compromise between the accuracy and speed of the measurements. When increasing the delay time, the relaxation seen in Fig. IV.9 also became less obvious, see Fig. IV.10. This was also the case for C2.

#### IV.3.2 Switching Currents

The switching currents for C2 and F1 were investigated as a function of the different in-plane magnetic field directions  $H_y, H_z$  after different field excursions. The switching currents were extracted by a thresholding algorithm. The offset voltage was calculated by computing the average of  $V$  in the range 500 nA to  $-500$  nA. This value was then subtracted from the voltage, and  $I_c^+$  ( $I_c^-$ ) was taken as the smallest (least negative) current for which the adjusted voltage was above (below)

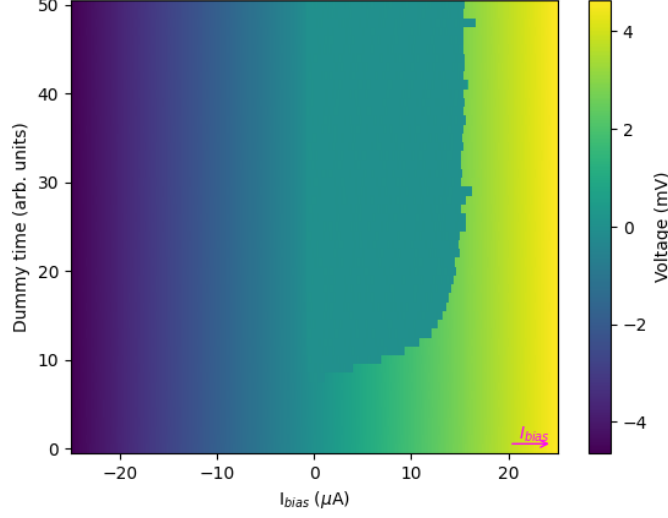


Figure IV.9: 50 I-V sweeps of C2, with the first sweep shown at  $y = 0$ . The current was swept from  $-25 \mu\text{A}$  to  $25 \mu\text{A}$ . Note that the switching current, visible as a sudden change in color at around  $15 \mu\text{A}$ , relaxes after  $\sim 20$  sweeps. Taken with a delay time of 0 s. Measured after a field excursion to  $\mu_0 H_y = 1 \text{ T}$ .

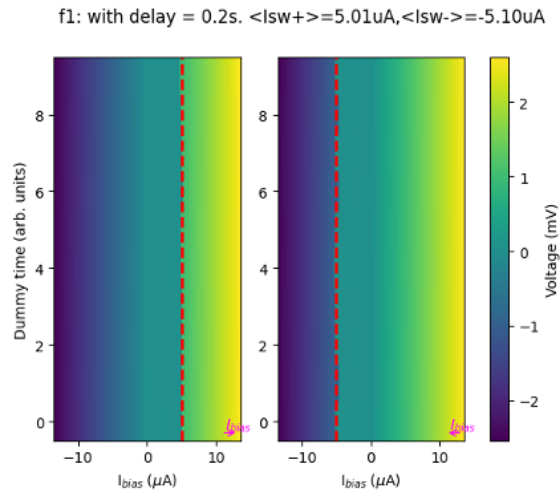


Figure IV.10: 10 IV-sweeps on F1 with 0.2 s delay time.  $y$ -axis denotes the sweep number, with the first sweep at  $y = 0$ . Average switching current indicated by red dotted line. Switching current was calculated for both current direction sweeps.

a predetermined tolerance. From trial and error, this tolerance was set to  $50 \mu\text{V}$ . During this section we will focus on the positive current sweep direction in order to show qualitative changes after different field excursions. The different current directions will be treated in the context of superconducting diode efficiency in Section IV.3.3. All measurements were performed with a delay time of 0.2 s as discussed in Section IV.3.1, and with a single IV-sweep at each point. If not stated otherwise, the magnetic field resolution was 2 mT. During these measurements it was found that:

- (i)  $I_c(H_y)$ ,  $I_c(H_z)$  have similar shapes.

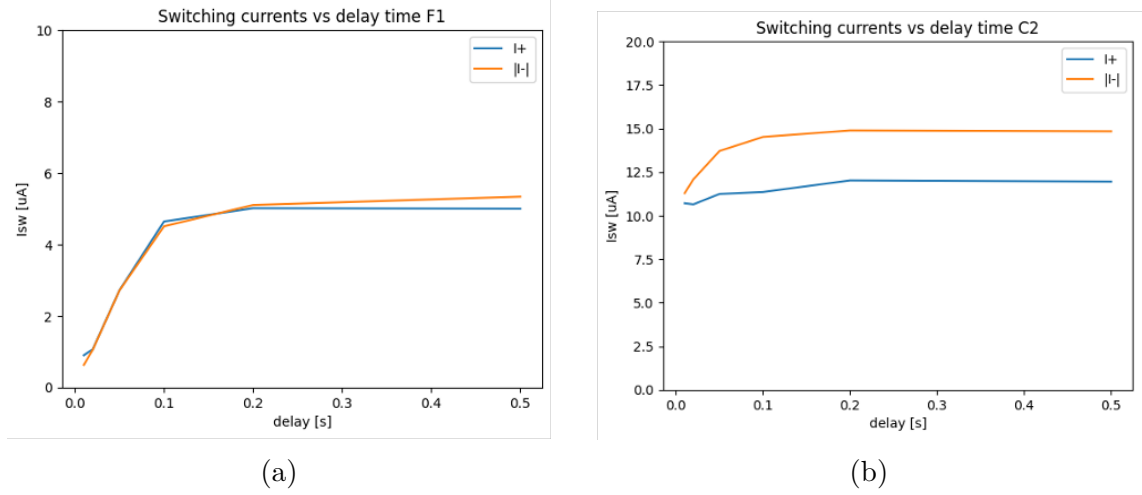


Figure IV.11: Average switching current as a function of delay time for both C2 and F1, for both up (blue) and down (orange) current sweeps. Each point taken from a plot such as Fig. IV.6. A strong dependence of the switching currents on the delay time can be seen, especially for F1.

- (ii) The devices exhibit strong hysteresis with regard to the magnetic field history.
- (iii) Both devices' switching currents have qualitatively similar dependence on the magnetic field history.

Because of observation (ii), the measurements were done by first performing a field excursion to some value in  $H_x, H_y$  or  $H_z$ , then ramping the field back to zero. After the excursion the switching currents  $I_c^+, I_c^-$  were measured. An example sequence to illustrate the procedure is:

1. Set  $\mu_0 H_z = 6$  T
2. Set  $\mu_0 H_z = 0$  T
3. Sweep  $\mu_0 H_y$  from 100 mT to  $-100$  mT, performing first one up current sweep followed by one current down sweep at each field value.
4. Sweep  $\mu_0 H_y$  from  $-100$  mT to 100 mT, performing first one up current sweep followed by one current down sweep at each field value.

The results from this procedure for F1 can be seen in Fig. IV.12a, where  $I_c^+$  for both the up and down field sweep directions is shown. As can be seen in this figure, the switching current as a function of field takes on a W-shape, something which was consistent after a large in-plane field excursion to  $H_z = 6$  T. This shape is also visible when measuring  $I_c(H_z)$  after a field excursion to  $H_z = 6$  T, as can be seen in Fig. IV.12b. A similar shape is visible for both devices measured, and for both current sweep directions. Since  $I_c(H_z)$  and  $I_c(H_y)$  showed qualitatively

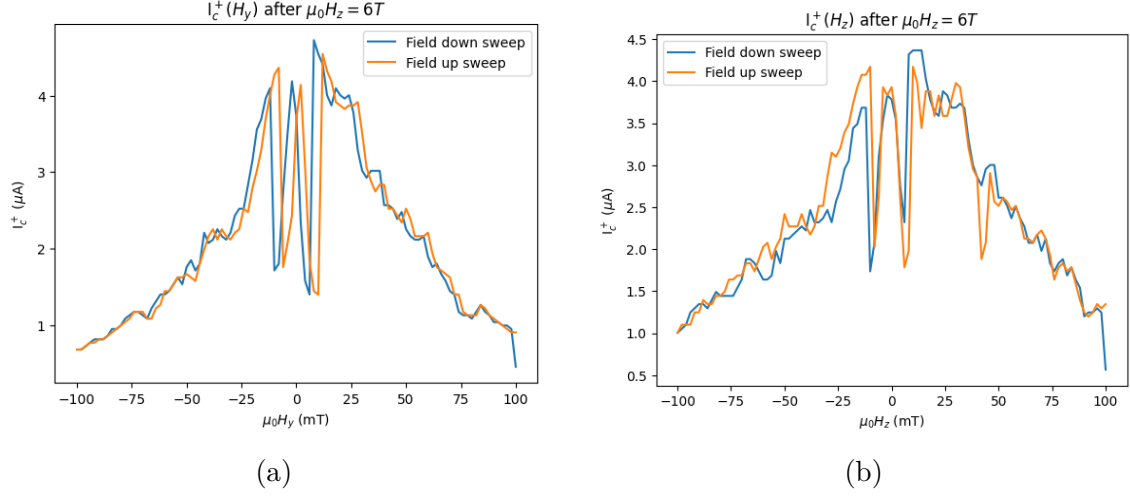


Figure IV.12: Illustration of the W-shape in the critical current after an excursion to  $\mu_0 H_z = 6$  T for F1. Shown are both (a)  $I_c^+(H_y)$  and (b)  $I_c^+(H_z)$ . In both figures a comparison between the up and down field sweep directions are shown. Some hysteresis is shown depending on field sweep direction, but qualitative features remain.

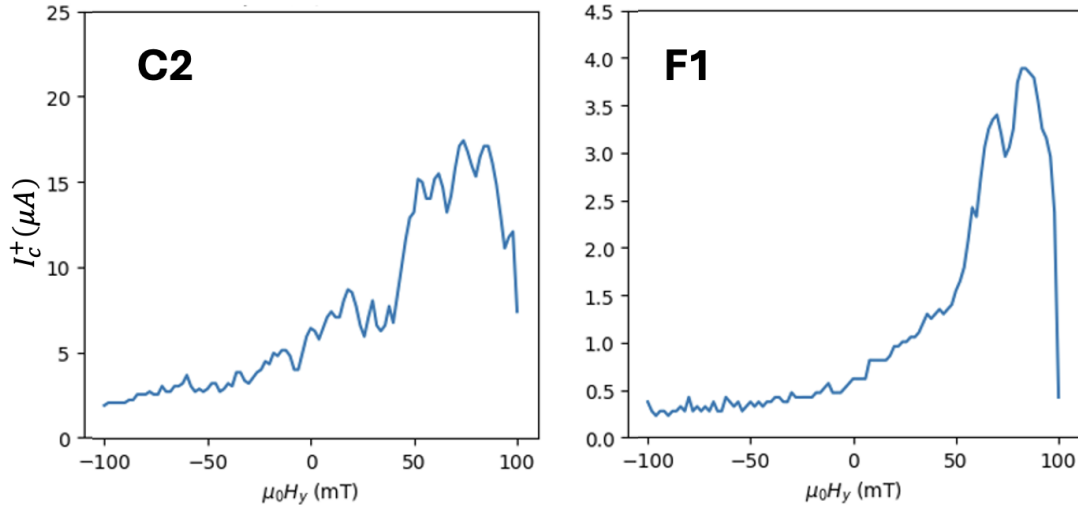


Figure IV.13: Switching current for C2 and F1 after out of plane field excursion to  $\mu_0 H_x = 1$  T. The peak is shifted to a positive value of  $\mu_0 H_y$ , compared to e.g. Fig. IV.12 where the main peak was centered around zero applied field.

similar behavior, it was decided to restrict measurements to  $I_c(H_y)$  for the rest of this section.

During measurements it was noted that the peak in  $I_c(H_y)$  could be shifted by performing out of plane field excursions, see Fig. IV.13. In this figure, taken after an out of plane field excursion to  $\mu_0 H_x = 1$  T, the peak in  $I_c(H_y)$  was shifted to some positive value of  $\mu_0 H_y$ . In order to investigate how this shift depended on the out of plane field strength, the following measurement sequence was set up:

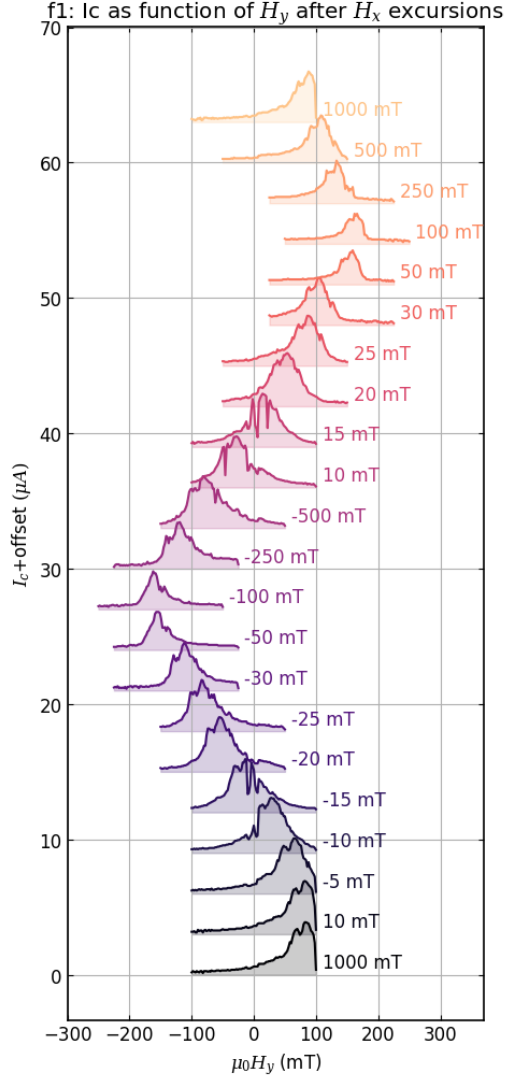
1. Perform a field excursion to  $\mu_0 H_z = 6$  T to get a known starting shape of  $I_c$ .

- (a) Set  $\mu_0 H_x = X$
- (b) Set  $\mu_0 H_x = 0$  T
- (c) Sweep  $\mu_0 H_y$  from 100 mT to  $-100$  mT with a step size of 2 mT and measure switching currents.
- (d) Repeat from Item 1a with the next value of  $X$ .

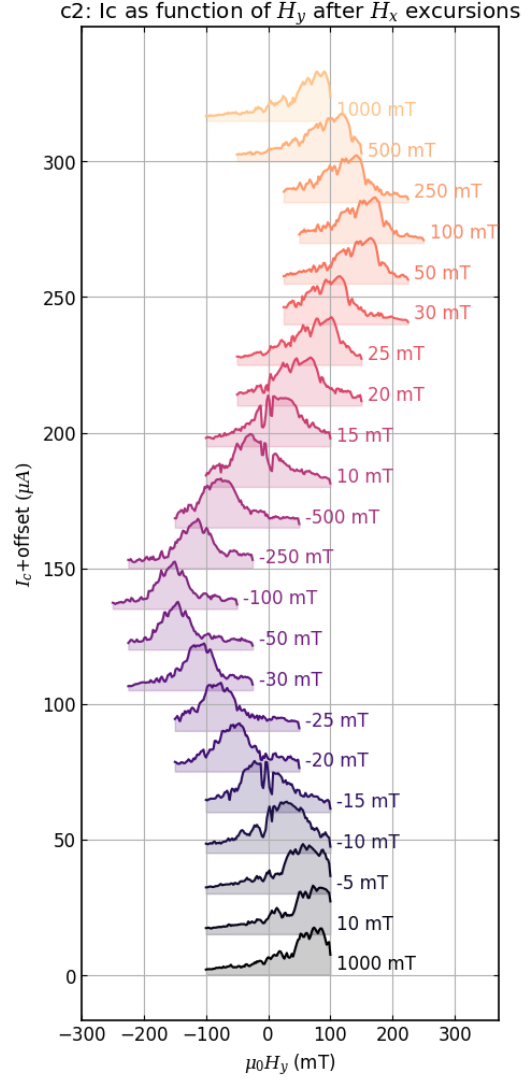
The results of these measurements can be seen in Fig. IV.14, where each subsequent curve is offset from the previous one by 3  $\mu$ A for F1 and 15  $\mu$ A for C2. This plot shows that the peak in critical current shifts to more negative  $H_y$  after negative  $H_x$  excursions, and to more positive  $H_y$  after positive  $H_x$  excursions. It can also be noted from the bottom two curves, see 1000 mT and 10 mT, that performing a smaller field excursion of the same polarity after performing a larger one does not change the position of the peak. However, after performing a large field excursion such as 1000 mT in the positive  $x$ -direction, it is enough to perform a small field excursion of the opposite polarity to change the position of the peak, see the  $-5$  mT excursion in Fig. IV.14. Consequently, a measurement protocol starting with a large field excursion of one polarity followed by successively larger field excursions of the other polarity is well justified. Furthermore, the W-shape in  $I_c(H_y)$  appears when the peak is close to  $H_y = 0$  T, see e.g the peaks labeled  $\pm 15$  mT in Fig. IV.14. This may be due to two different effects, one causing the W-shape around zero field and one moving the peak location. When the peak location happens to lie around zero these effects interfere and cause the W-shape.

Another surprising feature of Fig. IV.14 is the non-monotonicity of the peak shift with regard to the out of plane field excursions. This can be noted for both polarities of  $H_x$  and both devices, with a maximum shift after the  $|\mu_0 H_x| = 100$  mT excursion, and the peak shifting back toward zero for larger magnitude fields. This seems to indicate a magnetization that is higher for lower field values, which is something that was not expected.

To illustrate the dependence of the  $I_c$  peak position on the out of plane field excursion, a linear least squares fit to the function  $f(x) = Ae^{-B(x-x_0)^2} + C$  was performed on each of the  $I_c(H_y)$  plots seen in IV.14. The 10 mT measurement made directly after the 1 T excursion was excluded, see the discussion in the previous paragraph. Additionally, two measurements after excursions to  $-750$  mT and  $-1$  T were included. These were taken directly after the last 1 T excursion seen in IV.14. An illustration can be seen in Fig. A.2, where for each device one successful and one less successful fit is shown.  $\mu_0 H_y^{max} := x_0$  was then taken to be the peak offset and is plotted as a function of field excursions in Fig. IV.15. In this figure the non-monotonicity is clearly visible as the curve turns downward after  $\mu_0 H_x = 100$  mT and then possibly saturates at some offset between 50 mT and 100 mT. To see whether the shift saturates here a stronger magnet would be needed. The peak



(a)



(b)

Figure IV.14:  $I_c^+(H_y)$  after different out of plane field excursions for F1 and C2, field down sweep. The measured  $H_y$  range was adjusted to capture the peak after each excursion. First measurement is at the bottom, and subsequent curves are offset by (a)  $3 \mu\text{A}$  (b)  $15 \mu\text{A}$ . Next to each curve is noted the field excursion performed prior to that measurement.

shifts by a magnitude of  $\approx 300$  mT in  $\mu_0 H_y$  from the maximum positive shift to the maximum negative shift. The shift of the peak looks very similar for both devices.

That the peak is shifted in  $\mu_0 H_y$  indicates that the out of plane field excursion polarizes the CuMnAs, and that the magnetic field from this polarization cancels out the applied field at different values of  $\mu_0 H_y$  depending on the polarization. This is because the superconductor is expected to have the largest switching current at zero external field. As the magnetic moments closest to the interface in Fig. II.4 point in the same direction, it might be the case that these give rise to a local magnetic field near the interface, which could explain the offset in  $I_c(H_y)$ . However, an out of plane field is not expected to interact with the CuMnAs magnetic moments at all, since they lie in-plane. The non-monotonic behavior seen in IV.15 is also something that needs further investigation. That a smaller magnetic field induces a larger magnetization seems to suggest several magnetization scales. An example of a material that can have non-monotonic magnetization curves are ferrimagnets, which have opposing magnetic moments of different magnitudes, see e.g. Fig. 2b in Ref. [44]. As CuMnAs does not have magnetic moments of different magnitudes, the origin of this non-monotonicity remains unclear. As a normal metal near a superconductor also is made superconducting via the proximity effect [2], it might be the case that the CuMnAs/Al heterostructure near the interface is better modeled as an effectively unified superconductor. More research and modeling of the heterostructure is needed to understand the microscopic origins of these effects.

The W-shape seen in  $I_c(H_y), I_c(H_z)$  might be explained by spin canting and that the different spin sublattices of CuMnAs respond differently to applied fields. Spin canting is a phenomena where the different magnetic moments of an antiferromagnet are made to be not fully antiparallel. This may happen due to competition between the exchange interaction and the alignment of magnetic moments to external fields, which energetically favor antiparallel and parallel spin alignments, respectively. It can be imagined that applying a magnetic field at first tilts the magnetic moments of one sublattice, giving rise to a small net magnetization, resulting in the first dip in  $I_c(H)$  close to zero field. That  $I_c$  then rises again could be due to the other sublattice magnetic moments tilting to align in antiparallel with the first sublattice again, resulting in net zero magnetization and a larger switching current. This also assumes that one sublattice gets affected by the external field before the other.

During testing it was found that the shift of the peak in  $I_c(H_y)$  after  $H_x$  excursions could be reset by excursions to  $H_z = 6$  T, while excursions to the other in-plane direction  $H_y$  did not reset the structure. To determine whether this was due to the field strength in the  $y$ -direction being limited to 1 T, or whether it was due to the  $z$ -direction being more effective at resetting the structure, the following excursions were performed:  $\mu_0 H_x = -1$  T,  $\mu_0 H_y = 1$  T,  $\mu_0 H_z = 1$  T,  $\mu_0 H_z = 6$  T. The results can be seen in Fig. IV.16, and show that the structure does not reset

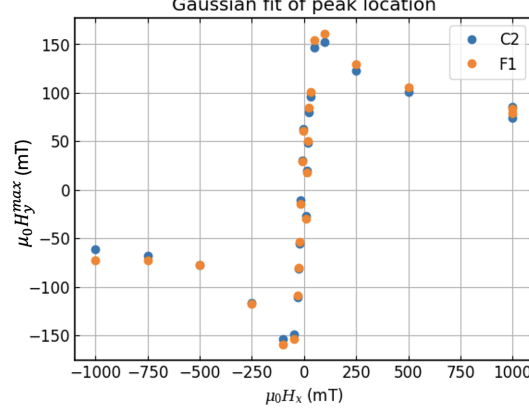


Figure IV.15: Location of peak of  $I_c$  in the  $I_c(H_y)$  curves as a function of out of plane field excursions for both devices, estimated by fitting to a gaussian peak. The non-monotonicity of the shift as a function of the excursions is clearly visible, with a maximum shift at around  $|\mu_0 H_x| = 100$  mT. Measurements taken in the same order as indicated in IV.14, with additional measurements at  $\mu_0 H_x = -750$  mT and  $\mu_0 H_x = -1$  T taken after the measurement sequence in IV.14.

fully after  $\mu_0 H_y = 1$  T, but seems to reset already after the  $\mu_0 H_z = 1$  T-excursion. That the shift could be reset by a large in-plane magnetic field further supports the claim that the shift is due to magnetization of the material.

### IV.3.3 Zero-Field Superconducting Diode Efficiency

During the exploration of the behavior of the switching currents, it was found that there was often a noticeable zero-field superconducting diode efficiency. An example is shown in Fig. IV.17, comparing  $I_c^+(H_z)$  to  $I_c^-(H_z)$  after a  $\mu_0 H_z = 6$  T field excursion. This measurement was taken simultaneously with the one shown in Fig. IV.12b. It can be noted that  $\eta$  is finite for a large portion of the range. As noted in Section II.1.2, the non-zero-field SDE could be attributed to an out of plane field component arising from a misalignment of the chip in the refrigerator with regard to the  $x$ -direction. Assuming a misalignment of  $2^\circ$  and an observable effect for an out of plane field of  $300 \mu\text{T}$ , this effect could explain the non-zero  $\eta$  for  $|\mu_0 H_z| \geq \frac{300 \mu\text{T}}{\sin 2^\circ} = 9$  mT. In Fig. IV.17 there is a finite diode efficiency at smaller fields than this, which must then be attributed to something other than chip misalignment.

Given the hysteretic behavior noted in Section IV.3.2 and the discussion accompanying Fig. IV.14 regarding polarization using a smaller magnetic field of the same polarity, the following procedure was decided upon, see also Fig. IV.18:

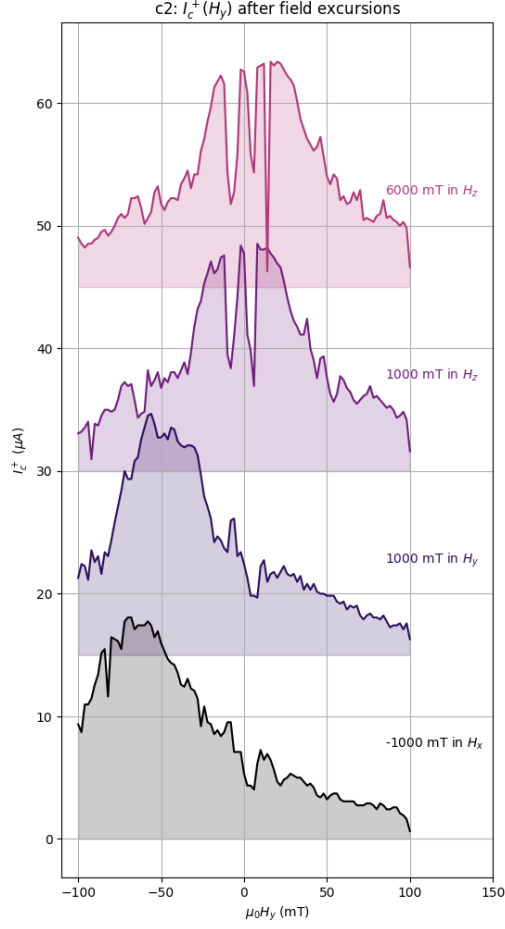


Figure IV.16:  $I_c^+(H_y)$  after field excursions to:  $\mu_0 H_x = -1 \text{ T}$ ,  $\mu_0 H_y = 1 \text{ T}$ ,  $\mu_0 H_z = 1 \text{ T}$ ,  $\mu_0 H_z = 6 \text{ T}$ .

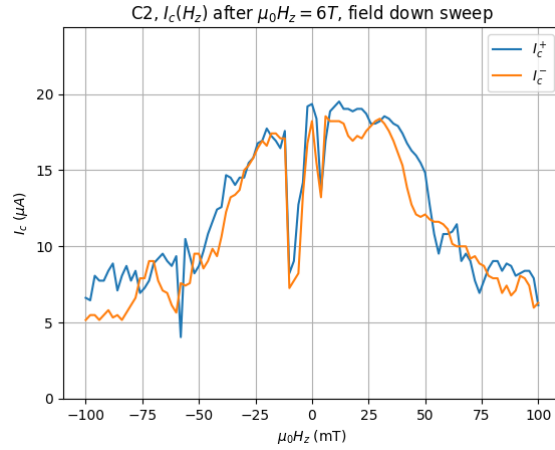


Figure IV.17:  $I_c^+(H_z), |I_c^-(H_z)|$  after  $\mu_0 H_z = 6 \text{ T}$  field excursion. Qualitatively similar results for up/down (blue/orange) current sweeps. Zero-field diode efficiency visible by the difference in switching currents at zero field, see Eq. (II.13).

1. Choose direction to polarize in. The directions investigated were  $x, y, z$  as well as  $45^\circ$  and  $135^\circ$  from the  $y$ -axis in the  $yz$ -plane. Example: choose  $H_y$ .

2. Set the field in the chosen direction to the maximum negative value. Example: set  $\mu_0 H_y = -1$  T.
3. Set the field in the chosen direction to zero. Example: set  $\mu_0 H_y = 0$  T.
4. Measure  $I_c^+$  and  $I_c^-$  ten times using 0.2 s delay time, discard the first two sweeps, and calculate  $\eta$  using  $\langle I_c^+ \rangle$  and  $\langle I_c^- \rangle$ . The first two sweeps are discarded to avoid any relaxation problems as the ones illustrated in Fig. IV.9.
5. Set the field in the chosen direction to the minimum positive value that is being investigated. Example: set  $\mu_0 H_y = 10$  mT.
6. Set the field in the chosen direction to zero. Example: set  $\mu_0 H_y = 0$  T.
7. Measure  $I_c^+$  and  $I_c^-$  ten times, discard the first two sweeps, and calculate  $\eta$  (Eq. (II.13)) using  $\langle I_c^+ \rangle$  and  $\langle I_c^- \rangle$ .
8. Repeat Item 5-Item 7 for the rest of the positive field values. Example: increase excursion field strength until  $\mu_0 H_y = 1$  T.
9. Repeat Item 5-Item 7 for the negative field values, starting with the smallest magnitude.

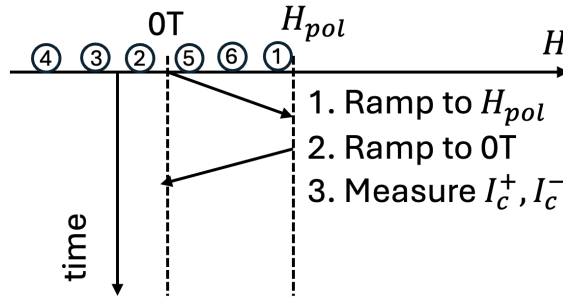


Figure IV.18: Illustration of the procedure used for measuring SDE. Material is polarized with some field  $H_{pol}$ , before  $\eta$  is measured at zero field. The circled numbers illustrate the order in which field excursions are taken.

The zero-field superconducting diode efficiencies as a function of prior field excursions are shown in Fig. IV.20 for C2 and in Fig. IV.21 for F1. Error bars show  $\pm 1\sigma_\eta$  as described in Section II.1.2, see also Eq. (II.18). In these figures a zero-field SDE of over 18 % is observed for both devices, the highest value occurring after the excursion to  $\mu_0 H_{yz} = -141$  mT with the field at a  $45^\circ$  angle from the positive  $y$ -axis. In Fig. IV.19 the positive and negative current sweeps of C2 after the excursion to  $\mu_0 H_y = -100$  mT are shown in the same plot, together with a histogram illustrating

the distribution of switching currents. This figure show a clear diode efficiency, and illustrates how  $\eta$  was calculated. Sign flips in the diode efficiency is seen in all cases for both devices, except for  $\eta(H_z)$  for F1. The diode efficiency at the end points only has different signs after the  $45^\circ$  excursion for F1, while it is seen for  $H_x$ ,  $H_y$ , and  $135^\circ$  excursions for C2.

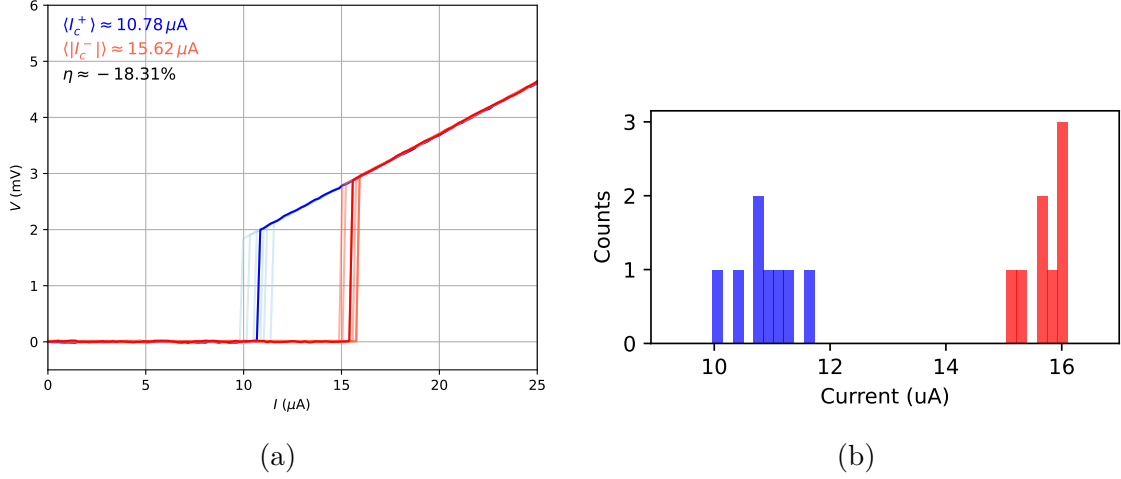


Figure IV.19: IV-sweeps for C2 at zero field after  $\mu_0 H_y = -100$  mT excursion. Shown are the eight different traces used in calculating the diode efficiency. (a) Raw data from upsweep together with  $|V(-I)|$  of down sweep, for ease of comparison. Bold lines are the traces closest to the average switching currents. (b) Histogram showing the distribution of the positive and negative switching currents of the traces in (a).

As seen in Fig. IV.20 and Fig. IV.21 the largest diode efficiencies were for most field directions at excursions to relatively smaller fields. The SDE measurements were only taken once, so further measurements might be needed to probe the overall structure of the diode efficiency curves.

That a zero-field diode efficiency is seen is also an indication that there is some local magnetization in the material. As mentioned previously, it might be due to the top layer of CuMnAs having aligned magnetic moments which could lead to a local magnetic field. Such a magnetization would break time reversal symmetry and could thus lead to a finite zero-field SDE. That the maximum of the efficiency is at comparatively smaller fields may be related to the non-monotonic magnetization that was speculated about in Section IV.3.2. That these phenomena might be related can be motivated by comparing Fig. IV.20b to Fig. IV.15, which show a remarkable likeness. However, if the out of plane magnetizes the CuMnAs, and the SDE is related to the magnetization, one might expect the largest  $\eta$  after the out of plane field excursions causing the largest magnetization, or shift in  $I_c(H_y)$ . This is not what is observed, as the maximum shift occurred after around  $\pm 100$  mT, while the maximum SDE after out of plane field excursions occurred at a magnitude of around 500 mT for both C2 and F1. The SDE was also larger after in-plane field

excursions, even though during the testing detailed in Section IV.3.2, these did not seem to polarize the material.

Another interesting thing to note is that the excursions in  $H_x$  seemed to, at first, suppress the switching currents, while higher field excursions increased the switching currents again. This happened for both devices, and is shown for C2 in figure Fig. IV.22. This is likely related to the behavior seen in Fig. IV.14b, but does not match exactly, likely because of the  $H_y$  field applied during these sweeps. The remaining switching currents used for calculating the diode efficiencies are shown in Fig. A.4 and Fig. A.3.

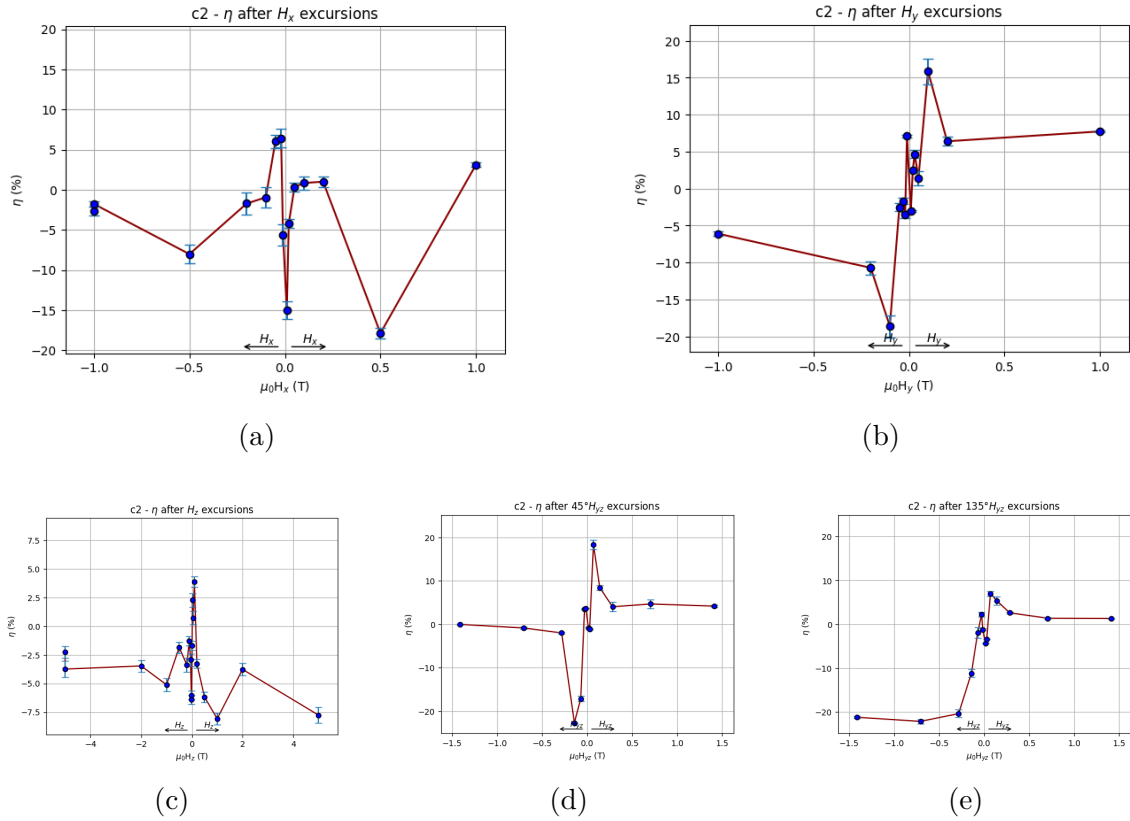


Figure IV.20: C2  $\eta$  at zero field as a function of different field excursions. (a)  $H_x$  excursions, (b)  $H_y$  excursions, (c)  $H_z$  excursions, (d)  $H_{yz}$  at  $45^\circ$  from positive  $y$ -axis excursions, (e)  $H_{yz}$  at  $135^\circ$  from positive  $y$ -axis excursions.

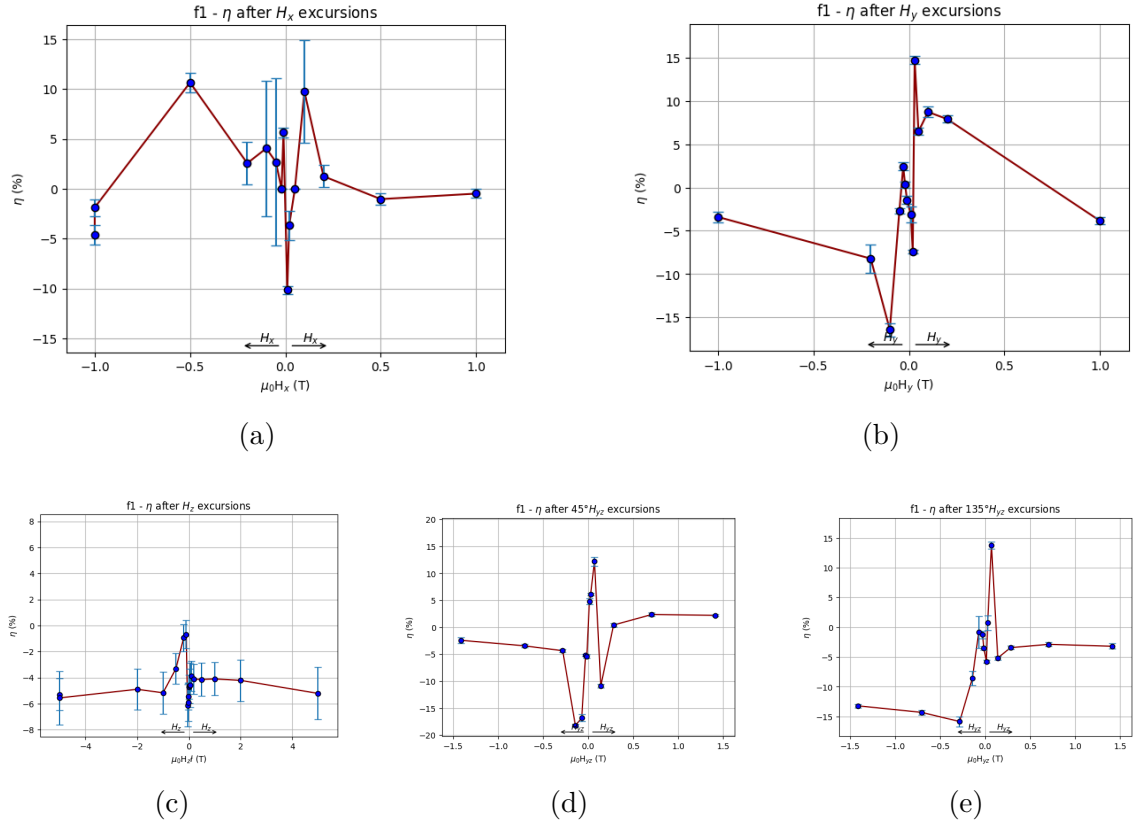


Figure IV.21:  $F1$   $\eta$  at zero field as a function of different field excursions. (a)  $H_x$  excursions, (b)  $H_y$  excursions, (c)  $H_z$  excursions, (d)  $H_{yz}$  at  $45^\circ$  from positive  $y$ -axis excursions, (e)  $H_{yz}$  at  $135^\circ$  from positive  $y$ -axis excursions.

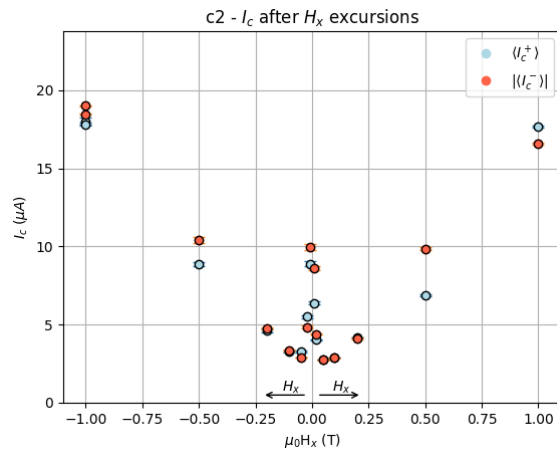


Figure IV.22: Switching currents for C2 at zero field as a function of out of plane field excursions. This data is the same as is used in Fig. IV.20a.

# Conclusion

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In this thesis the two magnetic material-superconductor heterostructures, MnTe/Al and CuMnAs/Al, have been examined. The MnTe/Al-based device showed a directional dependence of the sheet resistance, with  $R(T)$  increasing logarithmically with decreasing temperature. That no superconductivity was observed could be due to the magnetic moments in MnTe suppressing the superconductivity in Al. A possible explanation of the temperature dependence of the sheet resistance is proposed to be that the Mn formed an alloy with Al during growth, leading to the Kondo effect. Another possibility that would be consistent with measurements is weak localization leading to increased backscattering with lower temperatures. The directional dependence could either be due to differences in the crystallographic directions of MnTe, or due to variations in the fabrication. The next steps would be a structural analysis of the sample to check whether alloying has occurred, or to redo the measurements on a different device.

The CuMnAs/Al devices showed large hysteresis with regard to the field history. It also showed a W-shape in the switching currents as a function of the in-plane fields  $I_c(H_y)$  and  $I_c(H_z)$ . Furthermore, the peak in  $I_c(H_y)$  could be shifted by polarizing with an out of plane field excursion. The W-shape in the switching current as a function of in-plane field was speculated to arise from spin canting combined with different behavior as a function of external field of the two sublattices of CuMnAs. The shifted peak might be due to magnetization near the interface between CuMnAs and Al causing a local magnetic field. That this shift is non-monotonic, with a larger magnetization for comparatively smaller fields was unexpected. More research is needed to investigate the origin of this shift.

The hysteretic behavior of the CuMnAs/Al devices together with the observation that there seemed to be a zero-field SDE, led to exploring the zero-field SDE as a function of field excursions. As bulk CuMnAs is an antiferromagnet and thus does not produce any external magnetic fields, this zero-field SDE was not expected

initially. This is because a superconducting diode efficiency requires broken time-reversal symmetry, which is usually done by applying a magnetic field. During these measurements, the diode efficiency was indeed found to be dependent on the field history, and a maximum SDE of  $|\eta| \approx 23\%$  and  $|\eta| \approx 18\%$  was achieved for the two devices studied respectively. Like the shift in  $I_c(H_y)$  after  $H_x$  excursions, the zero-field diode efficiency might be attributed to uncompensated magnetic moments near the interface. These would lead to a local magnetic field which breaks  $\mathcal{T}$  symmetry, while the structure of CuMnAs breaks  $\mathcal{P}$  symmetry. Like the shift of  $I_c(H_y)$  after  $H_x$  excursions, the largest diode effects were found for comparatively smaller field excursions, leading to speculations that these effects might be related.

For future measurements of the CuMnAs/Al heterostructure, it would be desirable to obtain a higher resolution dependence of  $\eta$  on field excursions, in order to better understand how to optimize and tune the diode efficiency. It would also be of interest to retake the diode efficiency measurements, but "resetting" before each measurement by applying a  $\mu_0 H_z = 6\text{ T}$  field. This would be a way of ensuring that the field history is fully known. For example, the diode efficiency after  $H_y$  excursions to  $\pm 100\text{ mT}$  in Fig. IV.21 and Fig. IV.20 have a large value of  $\eta$ . However, a zero-field diode efficiency after  $\mu_0 H_y = 100\text{ mT}$  is not always observed, see e.g. Fig. IV.16. What can be seen in Fig. IV.16 is also that  $H_y$  does not drastically change the switching current structure, meaning that even the large field  $H_y$  excursions in Fig. IV.20 might be more dependent on what field excursions were performed prior.

To conclude, the MnTe/Al heterostructure has shown potential for directionally dependent resistance, but more measurements are needed to exclude the possibility of defects in the fabrication. More research is also needed to determine the origin of the increasing resistance with decreasing temperature. Furthermore, it has been shown that the CuMnAs/Al heterostructure is a promising material for exploring interactions between antiferromagnetism and superconductivity, and that it shows potential as a tuneable zero-field superconducting diode without stray field. More research is needed to fully characterize the behavior of this material and to understand the origin of the zero-field SDE and the effect of out of plane field excursions, especially the non-monotonic behavior of both these effects.

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# Appendix

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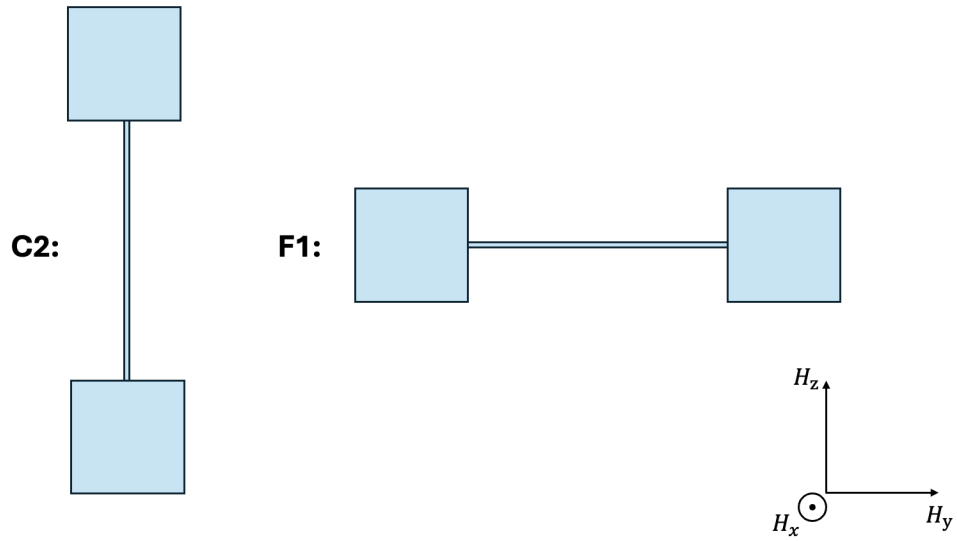
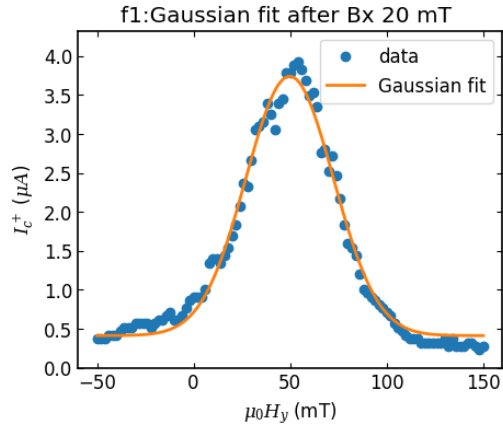
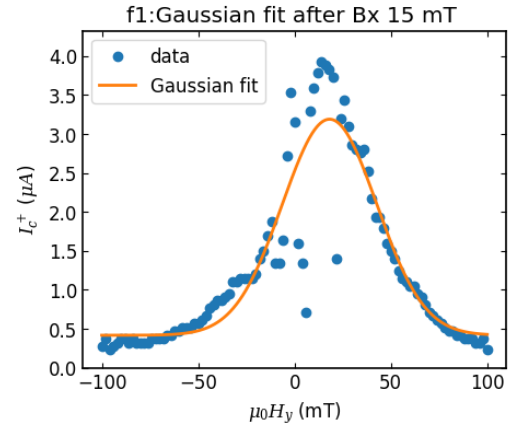


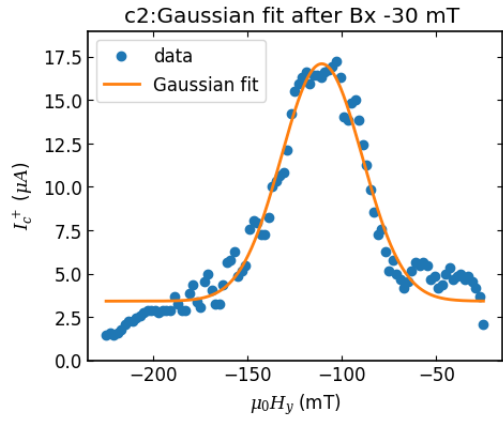
Figure A.1: Schematic of bar devices C2 and F1 illustrating their orientations with respect to the magnetic field directions  $H_x, H_y, H_z$ . For both devices  $L = 180\text{ }\mu\text{m}$ ,  $w = 5\text{ }\mu\text{m}$ .



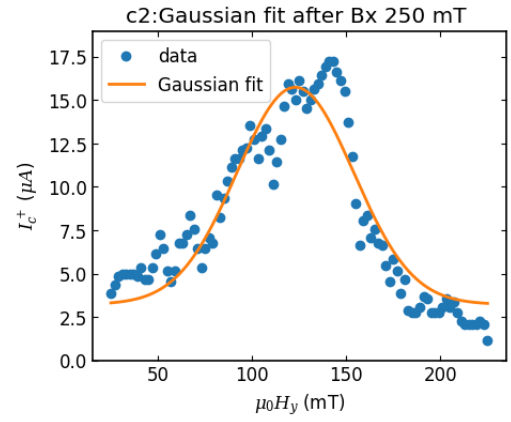
(a)



(b)



(c)



(d)

Figure A.2: Four examples of the Gaussian fits used to estimate the peak offset in  $I_c^+(H_y)$  for different out of plane field excursions. (a,b) For F1 (c,d) for C2. (a,c) More successful fits (b,d) Less successful fits.

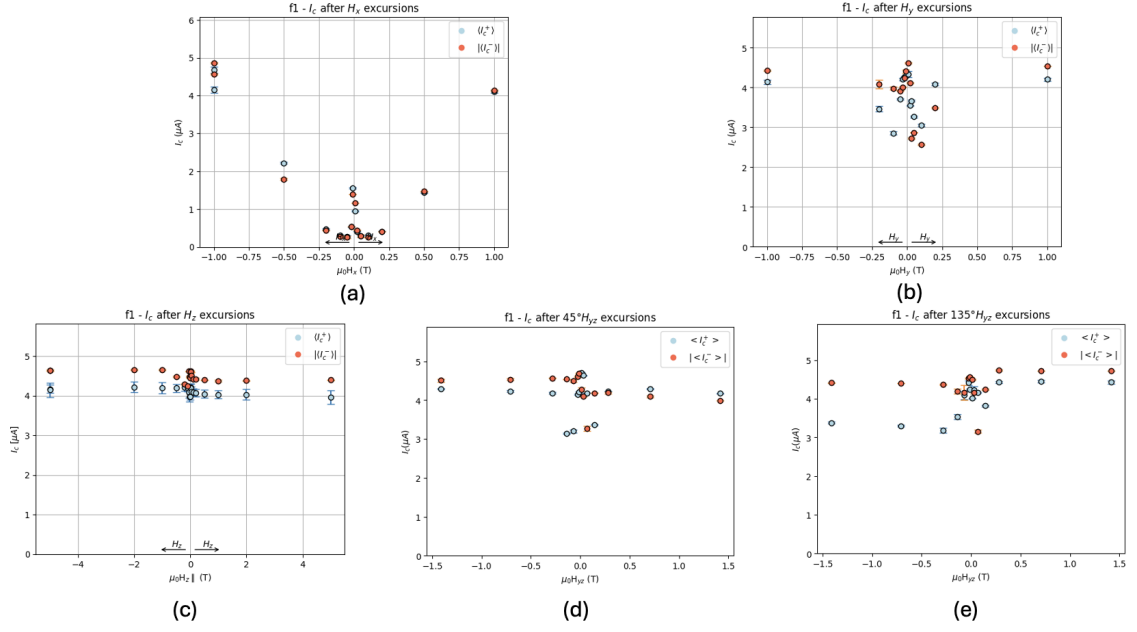


Figure A.3: Switching currents for F1 after (a)  $H_x$  excursions, (b)  $H_y$  excursions (c)  $H_z$  excursions, (d)  $H_{yz}$  at  $45^\circ$  from positive  $y$ -axis excursions, (e)  $H_{yz}$  at  $135^\circ$  from positive  $y$ -axis excursions.

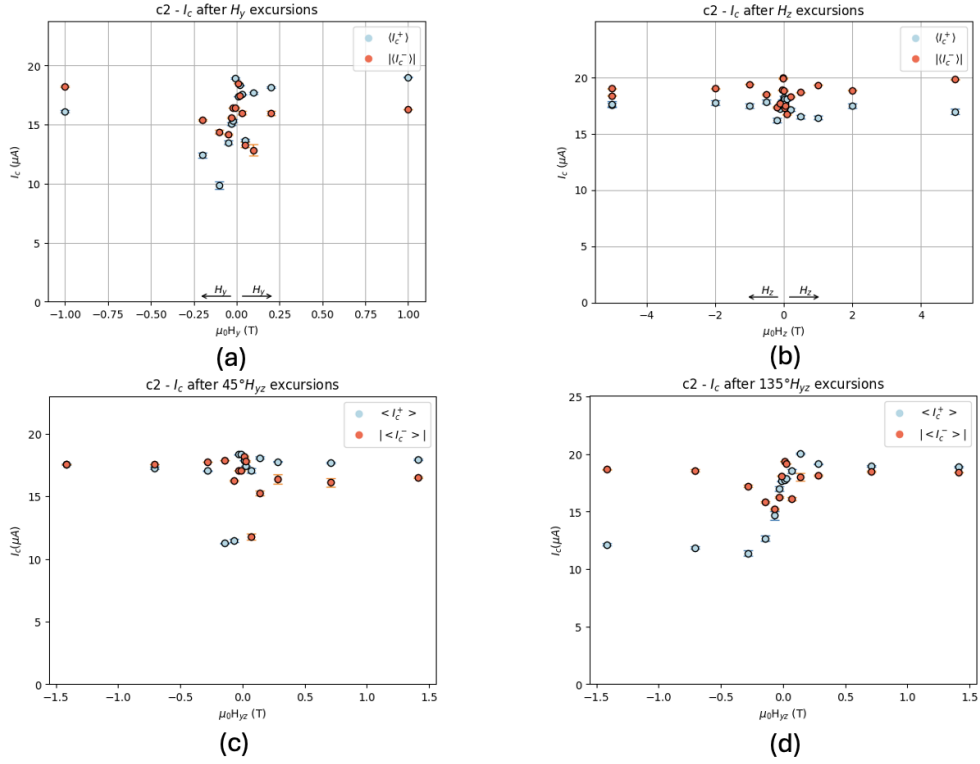


Figure A.4: Switching currents for C2 after (a)  $H_y$  excursions (b)  $H_z$  excursions, (c)  $H_{yz}$  at  $45^\circ$  from positive  $y$ -axis excursions, (d)  $H_{yz}$  at  $135^\circ$  from positive  $y$ -axis excursions. See Fig. IV.20a for the switching currents after out of plane excursions.