

HYDRA-TPC Prototype - a Time Projection Chamber for Light Hypernuclei Study at R^3B , GSI/FAIR

HYDRA-TPC-Prototyp – eine Zeitprojektionskammer für Studien zu leichten Hyperkernen bei R^3B , GSI/FAIR

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Darmstadt, Januar 20 2025

L. Ji

List of Figures

1.1. Pseudoscalar meson nonet as the function of Strangeness S , charge Q and third projection of isospin I_3	2
1.2. (a) Ground state baryon octet of spin 1/2 and (b) baryon decuplet of spin 3/2 as the function of Strangeness S , charge Q and third projection of isospin I_3	2
1.3. Λ decay widths in finite nuclei as a function of the mass number A . Γ_1 and Γ_2 denote one-body- and two-body- induced non-mesonic decay width; Γ_M and Γ_{NM} denote mesonic and non-mesonic decay widths. Γ_T is the total decay width. Reprinted figure with permission from [11] Copyright 2000 by the American Physical Society.	3
1.4. Average cross-section of $^{89}_{\Lambda}\text{Y}$ from (π^+, K^+) reaction, with the scattering angles ranging from 2° - 14° in the laboratory frame, as a function of Λ binding energy. F1 corresponds to Gaussian fitting peaks for main orbitals, and F2 corresponds to other bumps. Various lines correspond to different sets of values of the parameters including the potential depth, the diffuseness, and the volume integral of the potential. Reprinted figure with permission from [23] Copyright 2001 by the American Physical Society.	5
1.5. Neutron separation energies of Helium (hyper-)isotopes. Blue circles are from NCSM calculations (more details are in the following section) of helium isotopes, while green triangles and red squares are hypernuclear isotopes with cutoff momenta $p_\Lambda = 700 \text{ MeV}/c$ and $p_\Lambda = 600 \text{ MeV}/c$, respectively. Experimental values are shown in black bars or crosses for the (hyper-)nuclei. Vertical lines indicate extrapolation uncertainties. The figure is reprinted from [24] under CC BY 4.0.	6
1.6. Hierarchy of nuclear forces at increasing orders in ChEFT. Dashed lines represent pions and vertical solid lines are for nucleons. Filled circles show the coupling between nucleons and mesons, while filled squares stand for contact terms between nucleons. The $\nu = 1$ term does not exist due to the violation of parity conservation and time-reversal invariance. The figure is reproduced with permission from Springer Nature [27].	7
1.7. Calculation for the 1s, 1p, 1d, 1f, and 1g Λ single particle shells as a function of $A^{-2/3}$ where A is the mass number of the hypernucleus. The symbols represent experimental results. The figure is reproduced with permission from Springer Nature [28].	8
1.8. The EOS of SNM (blue) and PNM (black) as well as the symmetry energy (red) as a function of density obtained within the IU-FSU RMF model and the SHF approach using the SkIU-FSU parameter set. Reprinted figure with permission from [48] Copyright 2013 by the American Physical Society.	10
1.9. (a) The EOS pressure as a function of energy density with and without the presence of hyperons and (b) masses of neutron stars. A generic model with (black solid line) and without (red dashed line) hyperons has been considered. The horizontal lines show the observational mass of the Hulse–Taylor pulsar [51] and the recently observed PSR J1614-2230 [52] and PSR J0348+0432 [53]. Used with permission of The Royal Society (U.K.), from [54]; permission conveyed through Copyright Clearance Center, Inc.	11
1.10. Different stages of an event from HIC experiments.	14

1.11. Dimensionless mean square of radii and dimensionless two-body binding energy for halo systems. R represents the classical turning point and B represents the separation energy of the valence nucleon or hyperon. The filled circle denotes the deuteron, the filled squares denote nuclei where radii were extracted from experimental interaction cross-sections, the open squares are simple model estimates and the open circles are theoretical calculations. Used with permission of IOP Publishing, Ltd, from [118]; permission conveyed through Copyright Clearance Center, Inc.	17
1.12. Schematic drawing of the nucleus-nucleus collision in the framework of the Glauber model. The mass numbers of the projectile and the target are A and B . $\vec{b}$ is the vector for the impact parameter. $\vec{s}$ represents the distance between the center of the nucleus A and the interacting flux of nucleons. Used with permission of Annual Reviews, Inc., from [123]; permission conveyed through Copyright Clearance Center, Inc.	18
1.13. Total reaction cross-section of ^{22}C - ^{12}C collisions as a function of the rms radius of ^{22}C calculated by Monte Carlo Glauber theory. M represents the number of configurations. Reprinted figure with permission from [125] Copyright 2018 by the American Physical Society.	19
1.14. $^3_\Lambda\text{H}$ (left) lifetime, τ , and (right) binding energy, B_Λ , measurements obtained with different experimental techniques. The horizontal lines and boxes are the statistical and systematic uncertainties, respectively. The dashed and dash-dotted lines are the corresponding theoretical predictions. The figure is reprinted from [126] under CC BY 4.0.	20
1.15. Yield distribution of hypernuclei $^A_\Lambda\text{X}$ with different interaction cross-sections $\sigma_{\Lambda R}$. The ^{12}C beam intensity is 10^6 pps. The target thickness $d = 5$ cm (blue shaded band). The figure is reprinted from [138] under CC BY 4.0.	23
2.1. Mean energy loss rate of muons, pions, and protons in liquid (bubble chamber) hydrogen, gaseous helium, carbon, aluminum, iron, tin, and lead. The figure is reprinted from [144] under CC BY 4.0.	27
2.2. Number of primary ionization clusters for MIPs for various gases at the room temperature and the atmospheric pressure. Z_m and $\bar{Z}$ represent the total charge and the mean atomic charge of the gas molecule. The dashed line represents fit without the point of CF_4 . Reprinted from [148], with permission from Elsevier.	29
2.3. Ionization density as a function of the laser energy density. Three gas mixtures are (1) T2K (black): $\text{Ar}/\text{CF}_4/\text{iC}_4\text{H}_{10} = 95/3/2$, (2) P10 (red): $\text{Ar}/\text{CH}_4 = 90/10$, (3) Ar/CO_2 (blue) = $90/10$. Used with permission of IOP Publishing, Ltd, from [155]; permission conveyed through Copyright Clearance Center, Inc.	31
2.4. Number of ions collected on the electrodes as a function of applied voltage in a single wire gas chamber. The figure is reprinted from [156] under CC BY 3.0.	32
2.5. Relationships between absorption probability a and number of oxygen molecules N_{O_2} in the operating gas, $\text{Ar} + 3.45\% \text{CO}_2$, when the reduced electric field strength (the electric field strength over the gas pressure), $(E/P)_{cg}$, is (a) 107.3 V/cm/bar and (b) 214.5 V/cm/bar. $N_{\text{O}_2} = 10^{16} \text{ cm}^{-3}$ corresponds approximately to a fraction of 0.04% for oxygen and 0.23% for air in the gas. Reprinted from [160], with permission from Elsevier.	34
2.6. Electron drift velocity as a function of electric field strength in $\text{Ar} + \text{CO}_2$ mixtures at a pressure of 1 bar and temperature. Different markers indicate the fractions of CO_2 in the various gas mixtures. © 2006 IEEE. Reprinted, with permission, from [163].	35
2.7. Longitudinal and transverse diffusion constants D^* ($D^* = \sqrt{(2D/u)}$) of electrons as functions of the reduced electric field in different gases. Dash lines are calculated from Equation 2.21. The figure is reprinted from [165] under CC BY 4.0.	36

2.8. Paschen curves, the breakdown voltage as a function of the product of pressure and distance, for different gases. The figure is reprinted from [174] under CC BY 4.0.	40
2.9. Relation between discharge probability induced by a 15 GeV hadron beam and the detector gain on a Micromegas for various gas mixtures (more details about the Micromegas are in Section 2.3.2). $\langle Z \rangle$ denotes the mean atomic number of the gas mixtures. Reprinted from [175], with permission from Elsevier.	41
2.10. Electron collision cross-sections in (a) argon and (b) carbon dioxide. Reprinted from [177], with permission from Elsevier.	43
2.11. Schematic drawing of a time projection chamber. It usually consists of the drift cathode, the drift region, the amplification region, and the segmented anode.	44
2.12. Electric field inside holes of GEM. The induction field serves to extract electrons from the amplification towards the subsequent GEM or the anode. The figure is reprinted from [187] under CC BY 4.0.	45
2.13. Electric field inside holes of Micromegas with a micromesh-to-anode gap of 25 microns. The figure is reprinted from [195] under CC BY 3.0.	46
2.14. (a) Charge amplification and collection of the Micromegas detectors and (b) charge spreading on the DLC and induction on the pads for the resistive Micromegas detectors.	47
2.15. ENC versus input capacitance for different peaking times for a fixed gain. tp denotes the peaking time of the signal. © 2007 IEEE. Reprinted, with permission, from [206].	50
3.1. Schematic of the experimental setup. One event of ${}^3\text{H} \rightarrow \pi^- + p$ from the ${}^{12}\text{C} + {}^{12}\text{C}$ collision at 1.9 GeV/u is shown. The beam particle is shown in red, while the green curve is π^- and the blue curve is ${}^3\text{He}$. The figure is taken from [208].	51
3.2. Measurement of the magnetic field inside and in the proximity of GLAD. The color represents the ratio of the measured magnetic field to 1 Tesla, while the maximum current of the magnet corresponds to the maximum magnetic field of 2.2 Tesla. The figure is provided by Valerii Panin from GSI.	52
3.3. Position of HYDRA-TPC inside the funnel of GLAD.	53
3.4. Sectional view of the HYDRA-TPC prototype (schematic).	54
3.5. 3-Dimensional CAD sketch of the chamber with connectors and flanges.	55
3.6. (a) Electron drift velocity and (b) transverse and longitudinal diffusion coefficients for gas mixtures as a function of the electric field without the existence of the magnetic field. Calculation is based on Magboltz [213].	56
3.7. Top view of the field cage and its position in the chamber (schematic).	57
3.8. Wiring details on the PCB frame.	58
3.9. CAD design of the flexible PCBs.	58
3.10. (a) Assembly scheme of the field cage and (b) the pillars and the cathode-PCB.	59
3.11. CAD drawing for the top side of the GEM. Brown tracks denote conductive wires. The orange area denotes the GEM. The light green area denotes the PCB.	60
3.12. Cross section of the metal core pad plane (schematic).	61
3.13. CAD Drawing for the anode side of the metal core pad plane.	61
3.14. CAD Drawing for the front-end electronics side of the metal core pad plane.	61
3.15. Assembly of the field cage and the cathode on top of the bottom plate of the chamber.	62
3.16. Assembly of the HYDRA-TPC without windows. On top of the TPC are multiplexing boards, grounded with the chamber through yellow cables.	63

3.17. (a) Ratio of y component to the z component and (b) ratio of x component to the z component of the flux density in the active region when The TPC is positioned inside GLAD as Figure 3.3. The percentage is represented by variations in color. The xy plane is parallel to the cathode plane, while $z = 23$ cm is the distance from the cathode. The field map of GLAD was calculated by José Luis Rodríguez-Sánchez from Universidad de la Coruña.	64
3.18. (a) Schematic and (b) 3D CAD drawings of the mirror system. The system was built by Alexandru Enciu at TU Darmstadt. (c) A micro-mirror bundle, borrowed from the STAR collaboration. Each micro-mirror has a diameter of 1 mm perpendicular to the laser beam direction. The entire bundle has a diameter of 3 mm. Along the laser beam direction, the total length of the micro-mirrors is 5 mm. New micro-mirror bundles for the HYDRA-TPC are being constructed.	64
3.19. (a) An AsAd board connected with a ZAP board (employed to protect the AsAd board). (b) The ZC706 evaluation board with the z-CoBo firmware [218] developed by Mikołaj Ćwiok and Marcin Zaremba from the University of Warsaw.	65
3.20. (Top) Multiplexing boards installed on the pad plane. (Bottom) The scheme for one multiplexing board. The multiplexing boards are designed by Yelei Sun and constructed at the electronic workshop of TU Darmstadt.	67
4.1. Sectional view of the field cage (schematic). Axis z denotes the direction of the drift field. The black lines on the cathode denote the active area.	70
4.2. (a) Two-dimensional geometry input file for Gmsh and (b) result with the triangular meshing method. More details are in the text.	72
4.3. Equipotential lines calculated from ElmerFEM (Unit: V).	73
4.4. Ratio E_y/E_z at various positions. The horizontal position y ranges from 0 (at the center of the cathode) to 4.2 cm (at the edge of the active area). The vertical positions z are in the center of the drift field, 1 cm from the cathode and from the GEM respectively.	74
4.5. Ratio E_y/E_z for various construction tolerance of resistors. The horizontal position y ranges from 0 (at the center of the cathode) to 4.2 cm (at the edge of the active volume). The vertical position is fixed to $z = 1$ cm from the GEM.	75
4.6. Ratio E_y/E_z for a GEM-to-Micromegas distance of (a) 5 mm, with the bottom surface of the GEM located at $z = 32.00$ cm. and (b) 11 mm, with the bottom surface of the GEM located at $z = 31.40$ cm. The horizontal position y ranges from 0 (at the center of the GEM) to 4.2 cm (at the edge of the active volume). z denotes the distance from the cathode, as showcased in Figure 4.2.	76
4.7. The tracks of drift electrons starting at different y positions (from left to right, one at 0.5 cm from the inner layer of wires, one at the edge of the active volume, one inside the active area, two between the inner layer and the outer layer of wires). These electrons are emitted at $z = 5$ cm and are absorbed by the GEM or the wires. The black area is epoxy resin.	77
4.8. Distribution of Position deviation for 10000 electrons drifting 3 cm from the cathode at the edge of the active area.	78
4.9. Positional deviation of electrons in the drift region for different gas mixtures. The data points for each gas mixture, correspond to distinct y positions within the field cage (From left to right, $y = 4.2, 3.2, 2.2, 1.2$, and 0.2 cm). The default electric field is used. For both gas mixtures, the temperature is 293 K and the pressure is 1 bar.	79
4.10. Positional deviation of electrons for different d_t with the same electric field strength 1000 V/cm. For each point, there are 10000 electrons released at different y positions (From the left to right, $y = 0.2, 1.2, 2.2, 3.2$, and 4.2 cm).	80

5.1. Scheme for electrical connection and resistors of the HYDRA-TPC for gain tests. Red rectangles are wires and strips.	81
5.2. Measurement of the leakage current on the DLC with a bias of 440 V.	83
5.3. Setup for the gain tests with the ^{55}Fe source.	84
5.4. X-ray ADC spectrum on the micromesh. The test duration is 181 s. The electronics gain is $4\times$. The blue line is the raw spectrum while the red curve is the fitting result.	84
5.5. X-ray ADC spectrum on a group of pads. The test duration is 302 s. The amplification gain is $20\times$. The blue line is the raw spectrum while the red curve is the fitting result.	85
5.6. Relation between the total effective gain $G_{\text{tot_eff}}$ and the high voltage bias on the DLC layer U_{DLC} . The vertical error bars are smaller than the marker size and therefore not explicitly shown. The red line is the linear fit of the data points in the logarithmic scale.	86
5.7. Relation between the total effective gain $G_{\text{tot_eff}}$ and the transfer field E_t . The error bars are smaller than the marker size and therefore not explicitly shown. Lines of corresponding colors represent linear fits to the respective data points.	87
5.8. The relation between the total effective gain $G_{\text{tot_eff}}$ and the voltage applied on the GEM U_{GEM} . The error bars are smaller than the marker size and therefore not explicitly shown. Lines of corresponding colors represent linear fits to the respective data points in the logarithmic scale.	88
5.9. The relation between the total effective gain $G_{\text{tot_eff}}$ and the drift field E_d . The error bars are smaller than the marker size and therefore not explicitly shown.	89
6.1. Test measurement for different peaking times. For each peaking time, 500 laser events were taken from one AGET chip. Other common settings are: Writing clock frequency = 12.5 MHz; Trigger delay = 35 μs ; Range = 120 fC.	93
6.2. Sketch of the sectional view of the TPC and the laser system (not to scale). The black ellipses represent reflection mirrors. The gray ellipses represent the beam splitters. And the green ellipses are the bundles of the micro-mirrors. "Center" in the figure refers to the central position of the micro-mirrors. The values are taken from the CAD drawings for the HYDRA-TPC.	94
6.3. Raw waveform for 1000 events from 64 channels of one AGET chip. From the left to the right, the five signal peaks correspond to Peak 0 to 4. The electronics settings are: Peaking time = 232 ns; Writing clock frequency = 12.5 MHz; Trigger delay = 35 μs ; Range = 120 fC.	95
6.4. (a) Peak time and (b) measured drift velocity as a function of the drift field. Calculated drift velocities are from [213], based on Magboltz. The electronics settings are identical to those for Figure 6.3.	95
6.5. Temporal evolution of (a) peak amplitudes and (b) peak time for three laser peaks. Peak amplitudes and time were continuously extracted from the waveform for Channel 26 of a tested AGET chip with the same electronics configurations as in Figure 6.3.	96
6.6. (a) Sketch of the experimental setup for the first experiment. (b) Position of the HYDRA-TPC, the front-end electronics, and the reflection mirror in the GLAD magnet.	98
6.7. Electronics mapping in the pad plane for the first experiment. 'U' and 'L' denote the sockets connection for the multiplexing boards in Figure 3.20.	98
6.8. Raw waveform for each AGET chip of both CoBo 0 and CoBo 1 from 1000 events of Run 1 in the first experiment.	100
6.9. Temporal evolution of (a) peak amplitude and (b) peak time for three laser peaks in Channel 26, AGET 2, AsAd 0, CoBo 1 from Run 3 of the first experiment.	101

6.10. Temporal evolution of peak amplitudes for (a) Channel 17, (b) Channel 21, (c) Channel 30, and (d) Channel 34 of AGET 2, AsAd 0, CoBo 1 from Run 3 of the first experiment. These channels correspond to the pads with a relative X index -2, -1, +1, and +2 in Figure 6.7, which are neighboring to the pad for Channel 26 according to the multiplexing scheme in Figure 3.20.	101
6.11. Waveform of the 68 channels in AGET 2, AsAd 0, CoBo 0 of one event in Run 1 for the first experiment. (a) Raw waveform; (b) Waveform after FPN correction; (c) Waveform after noise subtraction.	102
6.12. Noise subtraction method. (a) Waveform of the noise and the signal channels; (b) Linear fitting of the relation between the noise and the signal in the time bin range [0,65]; (c) Waveform of the noise channel after scaling and shifting and the signal channel; (d) Waveform of the signal channel after correction.	102
6.13. Calculated standard deviation of the baseline in all AGET chips for (1) Run 1 and (2) Run 3. The black line represents the values before noise subtraction, while the red line represents the values after noise subtraction. The error bars are smaller than the marker sizes and therefore not explicitly shown.	103
6.14. All hits in the pad plane from one event for Peak 1 from Run 1 on CoBo 0 in the first experiment. The color denotes the signal amplitude in ADC channels above the baseline. The threshold is 90 in ADC channels above the baseline. Red dashed rectangles represent the area of each AGET chip detailed in Figure 6.7.	105
6.15. De-multiplexing method: (a) Selection of non-multiplexing pads and their neighboring multiplexed pads in the red rectangle; (b) Selection of the minimum-distance pads from the multiplexing group. The color of the boxes denotes the peak amplitude above the baseline in ADC channels. The red line is from the previous fitting. Red or yellow squares are from the same multiplexing group.	105
6.16. Track reconstruction in (a) the XY plane and (b) the XZ plane from all hits in Figure 6.14. The color of the boxes denotes the peak amplitude above the baseline in ADC channels. α denotes the angle between the reconstructed track and the XZ drift plane. β denotes the angle between the reconstructed track and the XY pad plane. The red solid lines represent linear fits to the hits.	106
6.17. Fitting parameters of the linear functions for all reconstructed tracks of Peak 1 from Run 1 on CoBo 0.	107
6.18. Angular distribution of the reconstructed tracks for (a) Peak 1 and (b) Peak 2 on CoBo 0 in Run 1. The angles are plotted in degrees.	108
6.19. Extrapolated Y positions of the micro-mirror from track reconstruction for (a) Peak 1 and (b) Peak 2 in Run 1 on CoBo 0.	108
6.20. Distances distribution of centroids of charge to reconstructed laser tracks in the pad plane for (a) laser Peak 1 and (b) Peak 2 from Run 1 in the first experiment.	109
6.21. Distances distribution of selected hits to reconstructed laser tracks in the drift plane for (a) laser Peak 1 and (b) Peak 2 from Run 1 in the first experiment.	109
6.22. Reconstructed tracks for laser Peak 1 from (a) Run 1 and (b) Run 3 of the first experiment. . .	111
6.23. Non-parallel component B_y of the magnetic field for (a) the horizontally placed TPC and (b) the tilted TPC at an angle θ	112
6.24. Laser alignment of the ^{12}C target and the TPC for the second experiment.	113
6.25. Electronics mapping for the second experiment. 'U' and 'L' denote the socket connection for the multiplexing boards in Figure 3.20.	113
6.26. Raw waveform for each AGET chip from all events of Run 13 in the second experiment. . . .	115
6.27. Raw waveform for each AGET chip from one event of Run 13 in the second experiment. . . .	116

List of Tables

3.1. ENC measurements of the GET system with the peaking time of 70 ns and 1014 ns separately.	66
6.1. Data for the first measurement campaign.	99
6.2. Fitting parameters and observables for Peak 1.	111
6.3. Fitting parameters and observables for Peak 2.	112
6.4. Comparison between the measured displacement and calculated values. B denotes the magnetic field in the GLAD magnet. Calculation 1 considers a positive magnetic field y component with a magnitude of 1% B . Calculation 2 considers a $\theta = 4.2^\circ$ tilting of the TPC in addition to Calculation 1. Calculation 3 considers a $\theta = 5.3^\circ$ tilting of the TPC in addition to Calculation 1.	112
6.5. Summary of data collected in the second experiment before accepting the ^{12}C beam. N represents the total number of events recorded by both CoBo 0 and CoBo 1.	114
6.6. Summary of distances between peaks for the first experiment, the second experiment, and from the design in Figure 6.2. '0 - 1' denotes the GEM-to-first bundle distance. '1 - 2' denotes the first-to-second bundle distance. '2 - 3' denotes the second-to-third bundle distance. '3 - 4' denotes the third bundle-to-cathode distance. Due to the absence of Peak 3 in the first experiment, the table displays the distance between Peaks 2 and 4.	115
A.1. Summary of data collected in February. 'L' represents the laser, 'B' the beam, 'LB' the simultaneous operation of laser and beam. 'LT' represents laser triggering, and 'PT' plastic wall triggering. The beam intensities in the table are expressed in units of particles per spill. 'TD' is short for trigger delay. Values marked with $\star$ were inferred from the counting rate of the LOS. The comment " I was unstable" indicates that during the measurement the beam intensity decreased to half of the maximum value at some point.	121

Zusammenfassung

Hyperkerne bieten einen einzigartigen Zugang zur Untersuchung von Hyperon-Nukleon-Wechselwirkungen. Ihre extrem kurzen Lebensdauern, in der Größenordnung von Sub-Nanosekunden, stellen jedoch erhebliche experimentelle Herausforderungen dar. Das HYpernuclei Decay at R³B Apparatus (HYDRA)-Experiment, das für den Betrieb innerhalb der R³B-Anlage bei GSI/FAIR konzipiert wurde, zielt darauf ab, Hyperkerne mit Hilfe von Schwerionenkollisionsexperimenten und hochpräziser invarianter Massenspektroskopie zu untersuchen. Diese Arbeit präsentiert die Entwicklung der HYDRA Zeitprojektionskammer (TPC), die speziell für die Spurverfolgung von π^- entwickelt wurde, die aus hypernuklearen Zerfällen innerhalb des GLAD-Magneten von R³B entstehen. Die TPC verfügt über einen zweischichtigen, drahtgestützten Driftfeldkäfig mit einer Driftstrecke von 300 mm und einer aktiven Fläche von $256 \times 88 \text{ mm}^2$. Eine hybride Verstärkungsstufe wurde implementiert, die aus einem Gas-Elektron-Multiplier (GEM) und einem Micromegas-Detektor besteht. Diese Konfiguration wird voraussichtlich einen Ionenrückfluss von weniger als 1% erreichen.

Das Design des Feldkäfigs wurde durch zweidimensionale Simulationen, unter Verwendung der Finite-Elemente-Methode und Monte-Carlo-Techniken optimiert, um ein homogenes Driftfeld sicherzustellen. Die Elektronendriftverschiebung wurde am Rand des aktiven Bereichs auf weniger als 250 μm und im zentralen Bereich der TPC auf weniger als 200 μm bestimmt.

Die Verstärkungsleistung der TPC wurde mit einer Röntgenquelle charakterisiert. Durch Anpassung der an die Elektroden angelegten Hochspannung, wurde der Einfluss variierender Hochspannungen in verschiedenen Bereichen auf die gesamte effektive Verstärkung der TPC quantifiziert.

Die TPC wurde erfolgreich in Betrieb genommen, wobei ein Front-End-Auslesesystem mit Multiplexeingängen und die digitalisierende Ausleseelektronik auf Basis des GET-Systems integriert wurden. Anschließend wurde die Spurverfolgungsleistung durch Messungen von Laser-Spuren bewertet, die von einer 266-nm-Ultraviolett-Laserquelle erzeugt und durch Mikrospeigel-Bündel, die in der TPC integriert waren, in das Driftvolumen reflektiert wurden. Ein Tracking-Algorithmus wurde für die Rekonstruktion dieser Laser-Spuren entwickelt. Experimentelle Ergebnisse zeigen eine räumliche Auflösung von besser als 3 mm in Driftrichtung, während die Pad-Plane-

Auflösung die gewünschte Anforderung von 200 μm nicht erfüllte. Zusätzlich wurde der Einfluss eines externen Magnetfelds auf die Driftbahnen der Elektronen innerhalb des GLAD-Magneten bei Magnetfeldstärken von 0 bis 0,92 T untersucht.

Abstract

Hypernuclei offer a unique approach to investigating hyperon-nucleon interactions. However, their extremely short lifetimes, on the order of sub-nanoseconds, pose significant experimental challenges. The HYpernuclei Decay at R³B Apparatus (HYDRA) experiment, designed for operation within the R³B setup at GSI/FAIR, aims to perform heavy-ion collision experiments with the primary objective of performing high-precision invariant mass spectroscopy of light hypernuclei.

This thesis presents the development of the HYDRA Time Projection Chamber (TPC) specifically designed for tracking π^- produced from hypernuclear decays within the GLAD magnet of the R³B. The TPC incorporates a double-layer wired drift field cage with a drift distance of 300 mm and an active area of $256 \times 88 \text{ mm}^2$. A hybrid amplification stage was implemented, comprising a Gas Electron Multiplier (GEM) and a Micromegas detector. This configuration is expected to achieve an ion back-flow of less than 1%.

The design of the field cage was optimized through two-dimensional simulations employing the finite element method and Monte Carlo techniques to ensure a homogeneous drift field. Electron drift displacement was determined to be less than 250 μm at the edge of the active region and less than 200 μm in the central region of the TPC.

The gain performance of the TPC was characterized using an X-ray source. By adjusting the high voltage applied to the electrodes, the influence of varying high voltages in different regions on the overall effective gain of the TPC was quantified.

The TPC was successfully commissioned with a front-end readout system incorporating multiplexing boards and digitizing readout electronics based on the GET system. Subsequently, its tracking performance was assessed through measurements of laser tracks generated by a 266-nm ultraviolet laser source and reflected into the drift volume by micromirror bundles, which were integrated within the TPC. A tracking algorithm was developed to reconstruct these laser tracks. Experimental results demonstrated a spatial resolution better than 3 mm in the drift direction, while the pad plane resolution did not meet the desired 200 μm requirement. Finally, the influence of magnetic fields on the drift electron trajectories was investigated within the GLAD magnet at magnetic field strengths ranging from 0 to 0.92 T.

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Contents

1. Introduction	1
1.1. Hypernuclei and Hyperon-Nucleon Interaction	1
1.1.1. Strangeness and Hyperons	1
1.1.2. Fundamental Concepts in Hypernuclear Physics	2
1.1.3. Nuclear Equation of State and Hyperons	9
1.2. Production of Λ -Hypernuclei in Experiments	12
1.2.1. Strangeness Exchange Reaction	12
1.2.2. Strangeness Production Reaction	13
1.2.3. Heavy Ion Collisions	13
1.3. Hypernuclear Halos and the Hypertriton Case	16
1.3.1. Halo Nucleus	16
1.3.2. Hypertriton	20
1.4. Introduction to HYDRA Experiment	21
1.4.1. Challenges and Solutions for Hypertriton Experiments	21
1.4.2. Experimental Method of HYDRA	22
2. Physical Background and Basic Concepts for Time Projection Chambers	25
2.1. Interaction of Radiation with Gases	25
2.1.1. Interactions of Heavy Charged Particles with Gases	25
2.1.2. Interactions of Low-energy Photons in the Gas	29
2.2. Working Principles of a Gaseous Detector	31
2.2.1. Types of Gaseous detectors	31
2.2.2. Transport of Electrons and Ions in Gas	33
2.2.3. Amplification	38
2.2.4. Discharges	39
2.2.5. Operation Gas for Gaseous Detectors	42
2.3. Gaseous Time Projection Chambers in Brief	43
2.3.1. Basic Concept	43
2.3.2. Micro Pattern Gaseous Detectors for Time Projection Chambers	45
2.3.3. Signal Processing	46
3. Design and Construction of HYDRA-TPC Prototype	51
3.1. Experimental Setup and the TPC	51
3.1.1. Overview	51
3.1.2. Technical Requirements for HYDRA-TPC Prototype	53
3.2. Main Components of HYDRA-TPC	55
3.2.1. Gas Chamber	55
3.2.2. Drift Field Cage and Cathode	57
3.2.3. Amplification Region and Pad Plane	59
3.2.4. Assembly of the TPC	62

3.3. Peripheral Equipment of HYDRA-TPC	63
3.3.1. Laser System	63
3.3.2. Signal Readout System	65
4. Drift Field Simulation	69
4.1. Simulation Methods and Software	69
4.2. Drift Field Simulation	71
4.2.1. Two-Dimensional Meshing	71
4.2.2. Drift Field Calculation	72
4.2.3. Drift field with different resistor chains	75
4.2.4. Calculation of the Transfer Field	75
4.2.5. Discussion	76
4.3. Electron Drift Simulation	77
4.3.1. Simulation Setup	77
4.3.2. Position Deviation of Drift Electrons	78
4.3.3. Drift Deviation in the Transfer Field	79
4.3.4. Discussion	80
5. Characterization of HYDRA-TPC	81
5.1. Test Preparation	81
5.1.1. Test Setup	81
5.1.2. Current Measurements	82
5.1.3. ADC Spectra of ^{55}Fe Source	83
5.2. Gain Studies	85
5.2.1. Gain Measurements	85
5.2.2. Discussion	89
6. Laser Measurements with HYDRA-TPC	91
6.1. Laboratory Measurements	91
6.1.1. Test Setup	91
6.1.2. Test Results	92
6.2. In-GLAD Measurements	97
6.2.1. Experimental Summary for the First Experiment	97
6.2.2. Waveform Analysis	99
6.2.3. Laser Track Reconstruction	104
6.2.4. Displacement of Tracks Under Magnetic Fields	110
6.2.5. Preliminary study for the Second Experiment	112
6.2.6. Discussion	116
7. Summary and Outlook	119
A. Data Overview for the February Measurement	121

1. Introduction

1.1. Hypernuclei and Hyperon-Nucleon Interaction

1.1.1. Strangeness and Hyperons

In 1950, the Λ^0 hyperon was first discovered in cloud chamber observations of cosmic rays conducted at an altitude of approximately 3 km [1]. Its distinctive characteristic was a neutral particle decaying into a V-shaped track pair observed within the chamber's photographs. It was also noted that its lifetime of approximately 10^{-10} seconds was significantly longer than the characteristic decay times of particles subject to strong interaction, typically around 10^{-23} seconds, or electromagnetic interactions, typically 10^{-16} - 10^{-21} seconds. To explain this anomaly, the concept of strangeness, denoted as S , was introduced as a conserved quantum number within the framework of strong and electromagnetic interactions [2][3]. The conservation of strangeness prevents the Λ^0 from decaying via these forces.

In the Standard Model (SM), there are six flavors of quarks and antiquarks, namely up, down, strange, charm, top, bottom quarks, and their corresponding antiquarks [4]. Nevertheless, only the strange quark (s) and strange antiquark ($\bar{s}$) can carry the strangeness quantum number of -1 and 1 , respectively. The strangeness [5] of a hadron is thus given by:

$$S = -(n_s - n_{\bar{s}}). \quad (1.1)$$

Strange quarks as well as up and down quarks are spin- $\frac{1}{2}$ fermions with electric charges of $-\frac{1}{3}e$, $+\frac{2}{3}e$, and $-\frac{1}{3}e$, respectively. The isospin third-projection I_3 for these quarks is 0 , $\frac{1}{2}$ and $-\frac{1}{2}$ correspondingly. I_3 is defined by

$$I_3 = \frac{1}{2}(n_u - n_d), \quad (1.2)$$

where n_u and n_d denote the quantity of up quark and down quark in the particle. The antiquarks possess the opposite charges and I_3 from the corresponding quarks. These quarks or antiquarks serve as fundamental constituents of hadrons. Within hadronic systems, the quarks carry color charges, designated as red, green, and blue, while antiquarks carry the corresponding anti-colors. In accordance with color confinement of the Standard Model [6], hadrons must be color neutral. Consequently, hadrons are restricted to specific combinations such as mesons, composed of a quark-antiquark pair with complementary color charges, and baryons, comprising color singlet states of three quarks, each possessing one of the colors.

In 1961, Murray Gell-Mann [7] and Yuval Ne'eman [8] independently proposed the eightfold way, a classification scheme for hadrons that established the foundation for the subsequent development of the quark model. This approach organized hadrons into SU(3) representation multiplets based on their quantum numbers, such as charge (Q), strangeness (S), and the third component of isospin (I^3). For mesons with spin-parity quantum numbers $J^P = 0^-$, known as pseudoscalar mesons, the meson nonet provides a classification scheme as illustrated in Figure 1.1. It includes mesons such as K^- with the quark content of $\bar{u}s$ and K^+ with the quark content of $u\bar{s}$ which possess non-zero charge and strangeness.

The baryon octet and decuplet, depicted in Figure 1.2, represent analogous classifications for baryons. Those

baryons which possess negative strangeness are called hyperons. There are in total four hyperon groups, namely Λ ($S = -1$), Σ ($S = -1$), Ξ ($S = -2$), and Ω ($S = -3$), distinguished by the strangeness and the third projection of isospin as showcased in Figure 1.2.

In their ground state, hyperons predominantly undergo weak decay rather than strong decay. This is primarily attributed to the conservation laws of energy, isospin, and strangeness within the framework of strong interactions. The weak decay process involves the conversion of a strange quark into an up quark accompanied by the emission of a W^- boson, accounting for the characteristic nanosecond lifetimes. The Σ^0 hyperon as the only exception, decays electromagnetically into a Λ^0 hyperon and a gamma ray. Other hyperons typically decay into nucleons or lower-strangeness hyperons. For instance, the Λ^0 hyperon, composed of uds quarks, has a mass of $1115.683 \pm 0.006 \text{ MeV}/c^2$ [9] and predominantly decays via $p^+ + \pi^-$ (63.9%) or $n + \pi^0$ (35.8%) [9] with a mean lifetime of approximately 261 ps [10].

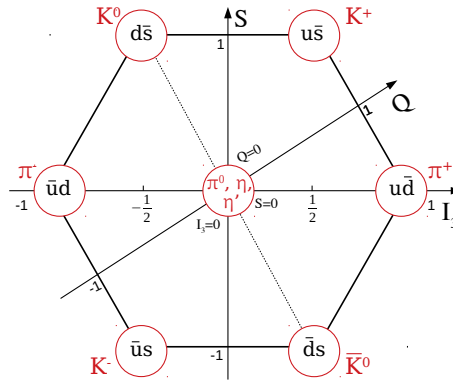


Figure 1.1.: Pseudoscalar meson nonet as the function of Strangeness S , charge Q and third projection of isospin I_3 .

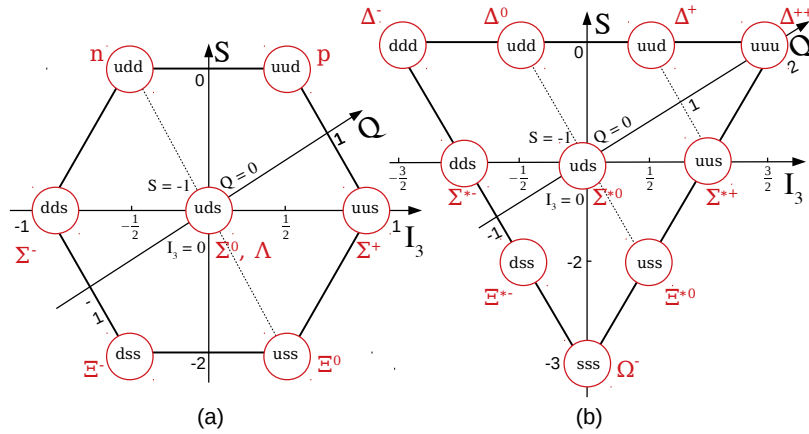


Figure 1.2.: (a) Ground state baryon octet of spin 1/2 and (b) baryon decuplet of spin 3/2 as the function of Strangeness S , charge Q and third projection of isospin I_3 .

1.1.2. Fundamental Concepts in Hypernuclear Physics

In 1952, the first evidence of a hypernucleus was observed in a photographic emulsion exposed to cosmic rays. An anomalous event was identified involving a secondary particle exhibiting the characteristic V-decay

topology of a Λ^0 hyperon. This suggested the existence of a nuclear system incorporating a bound Λ^0 particle, later termed as a hypernucleus. The hypernucleus is defined as a bound system composed of nucleons and minimally one hyperon. It is denoted as $^A_Y Z$, where A is the mass number; Z is the electric charge and Y indicates the hyperon(s). Thanks to the additional strangeness component, hypernuclei play a central role in the study of YN and YY interactions, where Y denotes hyperons and N denotes nucleons. The study of hypernuclei exhibits significantly more difficulty than normal nuclei due to their short lifetime. Since 1952, only approximately forty single- Λ , five double- Λ , and one single- Ξ hypernuclei have been identified.

Decay Modes of Hypernuclei

The strangeness $S = -1$ hypernucleus in the ground state can only decay via strangeness-changing weak interaction, though the mechanism differs from free hyperons due to the existence of nucleons. According to the Pauli exclusion principle, some momentum states are pre-occupied by the existing nucleons leading to the suppression of the mesonic decay of hyperons with the emission of low-momentum nucleons.

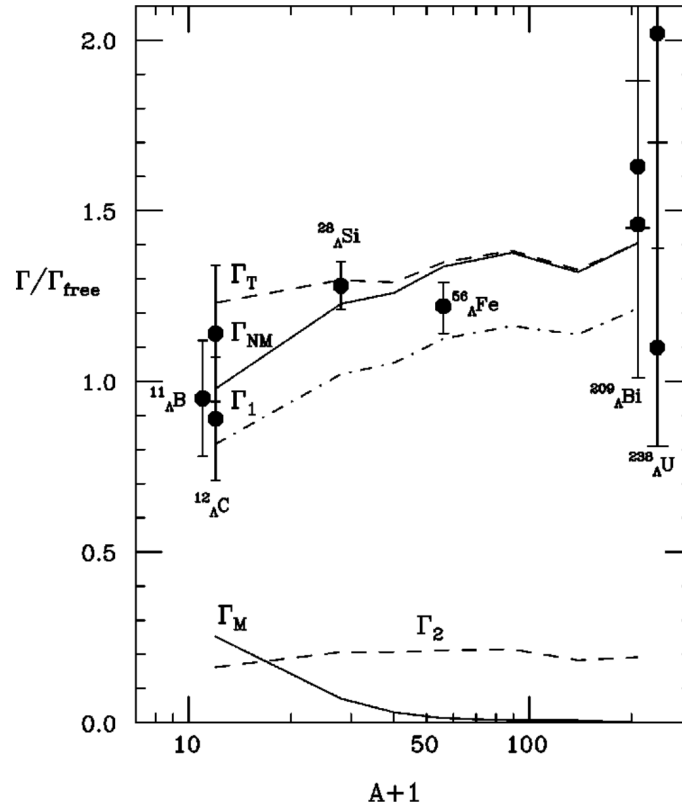


Figure 1.3.: Λ decay widths in finite nuclei as a function of the mass number A . Γ_1 and Γ_2 denote one-body- and two-body- induced non-mesonic decay width; Γ_M and Γ_{NM} denote mesonic and non-mesonic decay widths. Γ_T is the total decay width. Reprinted figure with permission from [11] Copyright 2000 by the American Physical Society.

For Λ hypernuclei, the Q -value, which is defined by the mass difference between the initial particles and the product particles in the reaction, is smaller than the Q -value for free Λ at rest (approximately $40 \text{ MeV}/c^2$ [12]). The momentum on the final nucleon after decay is roughly $100 \text{ MeV}/c$ [13], which is below the nucleon Fermi momentum $k_f \approx 250 \text{ MeV}/c$ given by Fermi Gas model:

$$k_f = \frac{\hbar}{R_0} \left(\frac{9\pi}{8} \right)^{\frac{1}{3}} = \frac{3\hbar}{4} \rho_0^{\frac{1}{3}}, \quad (1.3)$$

where $\rho_0 \approx 0.16 \text{ fm}^{-3}$ denotes the nuclear saturation density in a bound nucleus and $R_0 \approx 1.2 \text{ fm}$ denotes the approximate radius of a nucleon.

Although the suppression pronounces more for heavier hypernuclei (mass number $A > 5$) [14], the lighter hypernuclei remain to decay through mesonic decay modes since the hyperon being confined in a limited spatial region of a smaller hypernucleus allows for a wider momentum distribution while the local Fermi momentum at the surface is considerably smaller than the Fermi momentum in Equation 1.3 [15].

On the other hand, there is another decay mode for hypernuclei which is the non-mesonic weak decay (NMWD). In this mode, Processes like $\Lambda N \rightarrow NN$ and $\Lambda NN \rightarrow NNN$ take place without emission of pions. Instead, virtual pions or heavier mesons are produced during the weak decay of the hyperons and absorbed by nearby nucleons. Newly produced nucleons always have large momenta of approximately $400 \text{ MeV}/c$ for the one-nucleon-induced channel [16]. Consequently, this decay mode is not blocked by the Pauli principle and dominates over the mesonic decay. As presented in Figure 1.3, the mesonic decay rate decreases significantly with the increasing atomic number, while the non-mesonic decay channels dominate the $A > 10$ region.

For strangeness $S = -2$ hypernuclei, namely Ξ and double- Λ hypernuclei, new mechanisms are taken into account. The Ξ hyperon inside hypernuclei can directly decay via the strong force by $\Xi \rightarrow \Lambda\Lambda$ therefore preventing observation of the weak decay [17]. For double- Λ hypernuclei, nucleon-induced decay channels contribute dominantly while there is also a possibility of Λ induced decay channels [18].

The lifetime of hypernuclei is generally shorter than that of corresponding hyperons in free space due to NMWD with YN and YY interaction.

Structure of Λ -Hypernuclei

Thanks to its additional quantum number, a hyperon is a distinguishable particle from nucleons, thus it can occupy any orbital inside the nucleus without being blocked by the Pauli principle. Take the single- Λ hypernucleus ${}^A_\Lambda Z$ as an example, its Hamiltonian can be given by

$$H = H_N + T_\Lambda + \sum_{i=1}^{A-1} U_{i\Lambda} + \sum_{i>j=1}^{A-1} U_{ij\Lambda}, \quad (1.4)$$

$$H_N = \sum_{i=1}^{A-1} T_i + \sum_{i>j=1}^{A-1} V_{ij} + \sum_{i>j>k=1}^{A-1} V_{ijk}, \quad (1.5)$$

where T_i and T_Λ represent the kinetic energy operator for nucleon i and the Λ hyperon. V_{ij} , V_{ijk} , $U_{i\Lambda}$ and $U_{ij\Lambda}$ denote the NN, NNN, ΛN and ΛNN potentials respectively, which can be estimated with a phenomenological approach [19] [20]. As a simple approximation, the lambda hyperon can be assumed to be confined in a potential well generated by the nuclear core thus the Hamiltonian can be written as

$$H = H_0 + T_\Lambda + U_\Lambda, \quad (1.6)$$

where H_0 denotes the Hamiltonian of the nucleus core while U_Λ is the effective single particle central potential for the Λ hyperon whose major part can approximated using the Woods-Saxon potential:

$$U_\Lambda = V_0^\Lambda f(r) + V_{SO}^\Lambda \left(\frac{\hbar}{mc} \right)^2 \frac{1}{r} \frac{df(r)}{dr} (\vec{l} \cdot \vec{\sigma}), \quad (1.7)$$

$$f(r) = \{1 + \exp[(r - R)/a]\}^{-1}, \quad (1.8)$$

where V_0^Λ represents the potential depth, V_{SO}^Λ represents the spin-orbit potential, a denotes the diffuseness and R is the nuclear radius. The first term on the right side of Equation 1.7 is the common Woods-Saxon potential and the second term accounts for the coupling between the hyperon's spin and its orbital angular momentum. The spin-orbit interaction of ΛN in a hypernucleus is considerably smaller than the NN interaction [21], thus the general values of these parameters for the Λ hyperon [22] are given by:

$$\begin{aligned} R_\Lambda &= 1.1(A - 1)^{1/3} \text{ fm}, \\ a_\Lambda &= 0.60 \text{ fm}, \\ V_0^\Lambda &= 30 \text{ MeV}, \\ V_{SO}^\Lambda &= 2.0 \text{ MeV}. \end{aligned}$$

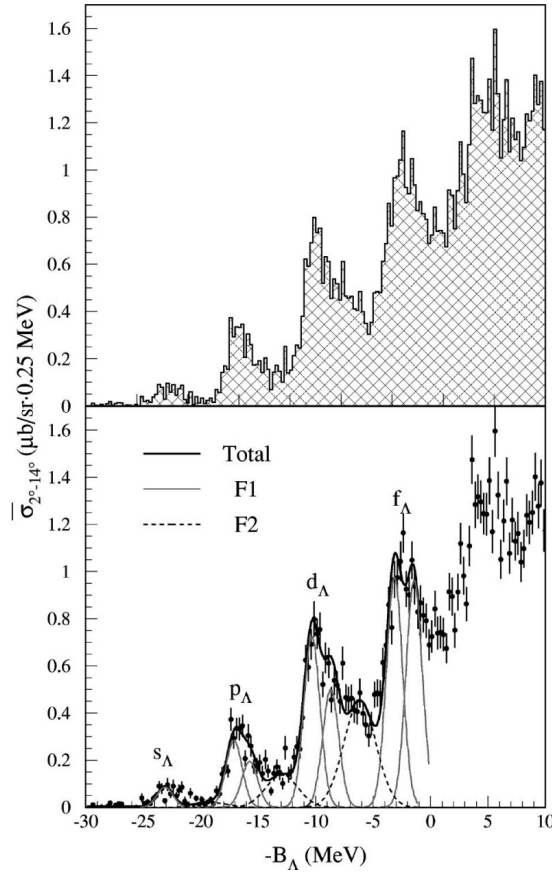


Figure 1.4.: Average cross-section of $^{89}_\Lambda\text{Y}$ from (π^+, K^+) reaction, with the scattering angles ranging from 2° - 14° in the laboratory frame, as a function of Λ binding energy. F1 corresponds to Gaussian fitting peaks for main orbitals, and F2 corresponds to other bumps. Various lines correspond to different sets of values of the parameters including the potential depth, the diffuseness, and the volume integral of the potential. Reprinted figure with permission from [23] Copyright 2001 by the American Physical Society.

Experimental investigations of hypernuclei typically involve the extraction of observables, such as the binding

energy of the Λ hyperon in a hypernucleus, B_Λ , defined by

$$B_\Lambda({}^A_\Lambda Z) = B_{tot}({}^A_\Lambda Z) - B_{tot}({}^{A-1}Z). \quad (1.9)$$

where B_{tot} denotes the total binding energy of a nucleus or hypernucleus. For a specific hypernucleus, the potential parameters are adjusted to reproduce experimental binding energy values. For instance, a spin-orbital potential of $V_{SO} = 4$ MeV was obtained from the experimental mass spectrum of ${}^{89}_\Lambda\text{Y}$ with a small spin-orbital splitting, as illustrated in Figure 1.4. Measuring binding energies is crucial for phenomenological studies to acquire more precise parameter values for potential calculation.

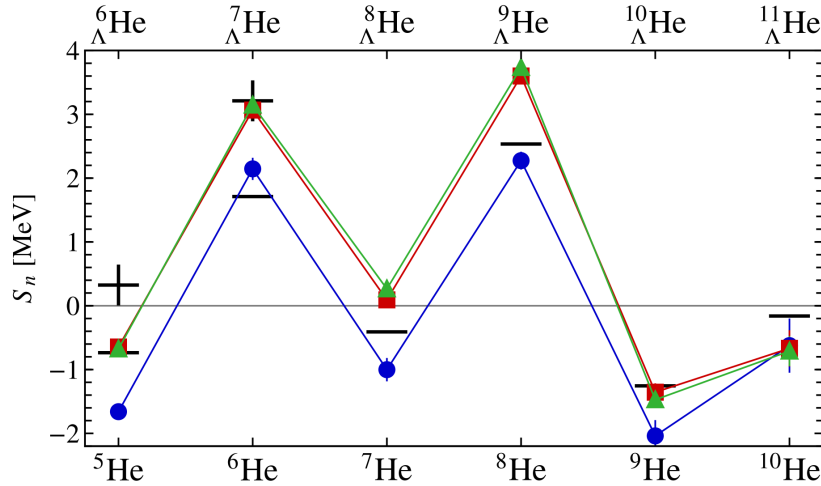


Figure 1.5.: Neutron separation energies of Helium (hyper-)isotopes. Blue circles are from NCSM calculations (more details are in the following section) of helium isotopes, while green triangles and red squares are hypernuclear isotopes with cutoff momenta $p_\Lambda = 700$ MeV/ c and $p_\Lambda = 600$ MeV/ c , respectively. Experimental values are shown in black bars or crosses for the (hyper-)nuclei. Vertical lines indicate extrapolation uncertainties. The figure is reprinted from [24] under CC BY 4.0.

Λ hyperons introduce an additional attractive force within hypernuclear systems, exhibiting an enhanced binding energy which results in a shift of the neutron and proton drip lines, defined as the last bound neutron-rich or proton-rich nuclei, from their normal limits. For instance, ${}^5\text{H}$ and ${}^6\text{Be}$ are unbound nuclei which are beyond the neutron and proton drip lines respectively, while their corresponding Λ hypernuclei, ${}^6_\Lambda\text{H}$ is bound [25] and ${}^7_\Lambda\text{Be}$ is theoretically bound [26] though experimentally not yet observed.

The separation energy of removing a proton S_p or neutron S_n from nucleus ${}^A Z$ is defined by

$$\begin{aligned} S_n &= M({}^A Z) - M({}^{A-1}Z) - M_n \\ &= B_{tot}({}^A Z) - B_{tot}({}^{A-1}Z), \end{aligned} \quad (1.10)$$

$$\begin{aligned} S_p &= M({}^A Z) - M({}^{A-1}(Z-1)) - M_p \\ &= B_{tot}({}^A Z) - B_{tot}({}^{A-1}(Z-1)). \end{aligned} \quad (1.11)$$

The glue-like feature of Λ hyperons is evidenced when

$$S_p({}^{A+1}_\Lambda Z) > S_p({}^A Z),$$

at the proton-drip line, or

$$S_n({}^{A+1}_\Lambda Z) > S_n({}^A Z),$$

at the neutron-drip line. Calculations of neutron separation energies for helium (hyper-)isotopes demonstrate this, as illustrated in Figure 1.5. Moreover, the negative neutron separation energy of ${}^7_\Lambda\text{He}$ and the positive neutron separation energy of ${}^8_\Lambda\text{He}$ demonstrate the potential of Λ hyperons to extend the neutron drip line beyond that of the corresponding nuclear isotopes.

YN Interactions

While phenomenological models including the Woods-Saxon potential provide valuable insights, they often cannot offer a systematic theory for precise calculations and a deeper understanding of YN and YNN interactions. These baryon-baryon interactions represent the residual effects of the strong force. While Quantum Chromodynamics (QCD) provides the fundamental theory of the strong force, its direct application to baryon-baryon interactions is challenging due to the non-perturbative nature of the interactions at low energies when quarks are confined inside the hadrons. Chiral Effective Field Theory (ChEFT) offers a powerful framework for studying baryon-baryon interactions at this energy range. This approach exploits the approximate chiral symmetry of QCD to construct an effective theory that focuses on hadrons, such as nucleons and pions, as the fundamental degrees of freedom, rather than quarks. By considering the exchange of mesons, particularly pions, ChEFT provides an efficient way to describe baryon-baryon interactions.

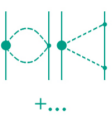
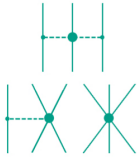
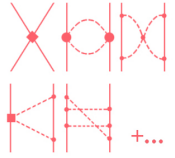
	2N Force	3N Force	4N Force
LO (Q/Λ) ⁰			
NLO (Q/Λ) ²			
NNLO (Q/Λ) ³			
N³LO (Q/Λ) ⁴			

Figure 1.6.: Hierarchy of nuclear forces at increasing orders in ChEFT. Dashed lines represent pions and vertical solid lines are for nucleons. Filled circles show the coupling between nucleons and mesons, while filled squares stand for contact terms between nucleons. The $\nu = 1$ term does not exist due to the violation of parity conservation and time-reversal invariance. The figure is reproduced with permission from Springer Nature [27].

The potentials in Equations 1.4 and 1.5 can be systematically expanded in power terms of the small external

momentum scale of a nucleon or the pion mass P over the large chiral symmetry breaking scale $\Lambda_\chi \sim 1$ GeV. The expansion is organized in the form of $(P/\Lambda_\chi)^\nu$, with the non-negative integer ν , which determines the power counting scheme as illustrated in Figure 1.6. In the lowest order, the leading order (LO), there are the momentum-independent contact term as the four-leg-nucleon diagram in the first row of Figure 1.6 and the one-pion exchange (OPE) term as the second diagram. In higher orders, additional contributions become significant, such as derivatives of the contact term and multi-pion exchange terms. To avoid divergence of the expansion, arising from factors such as short-distance singularities and high-energy particles, minimum-distance cutoff or maximum momentum cutoff is usually applied. As a consequence, the precision of ChEFT calculation results is sensitive to the selection of the cutoff scale and the inclusion of specific interaction diagrams.

The expansion coefficient for each power term is called the low-energy constant (LEC). These LECs control the short-range behavior of ChEFT and are not directly calculated from ChEFT but must be fitted to experimental data or calculated from QCD theories. ChEFT has been successfully applied to nucleonic two-body interactions up to next-to-next-to-next-to-next-to-leading order (N^4 LO) [29] and three-body interactions up to next-to-next-to-next-to-leading order (N^3 LO) [30].

While ChEFT has proven a practical and accurate approach to describing NN interactions, its extension to YN interactions presents unique challenges. The distinct nature of YN interactions, characterized by the distinguishability of hyperons and nucleons, allows for antisymmetric spin-orbit terms, such as $(\vec{s}_1 - \vec{s}_2) \cdot \vec{L}$, that are forbidden in NN interactions. The YN interaction is generally weaker than the NN interaction, demanding a tailored approach. Moreover, the small mass difference of ~ 80 MeV/ c^2 between Λ and Σ hyperons introduces the possibility of Λ - Σ conversion, which requires the inclusion of Λ N- Σ N coupling terms in YN one-boson-exchange (OBE) models [31]. However, these coupling terms remain a source of significant ambiguity. The scarcity of experimental YN data has restricted systematic investigations of ChEFT to LO and next-to-leading-order (NLO) calculations which are only applicable to two-body YN interactions. A comparison of calculation results employing YN interactions derived at LO [32], 2013 (NLO13 [33]) and 2019 (NLO19 [34]) versions of NLO together with experimental data is drawn in Figure 1.7. The LO calculation overestimated the attraction of YN interaction and NLO13 underestimated it, while NLO19 is best consistent with the data among the three models yet it fails to predict the data points for $^{13}_\Lambda\text{C}$ and $^{208}_\Lambda\text{Pb}$. More YN scattering experiments are necessary to provide better constraints on the models and reduce uncertainties for higher-order ChEFT

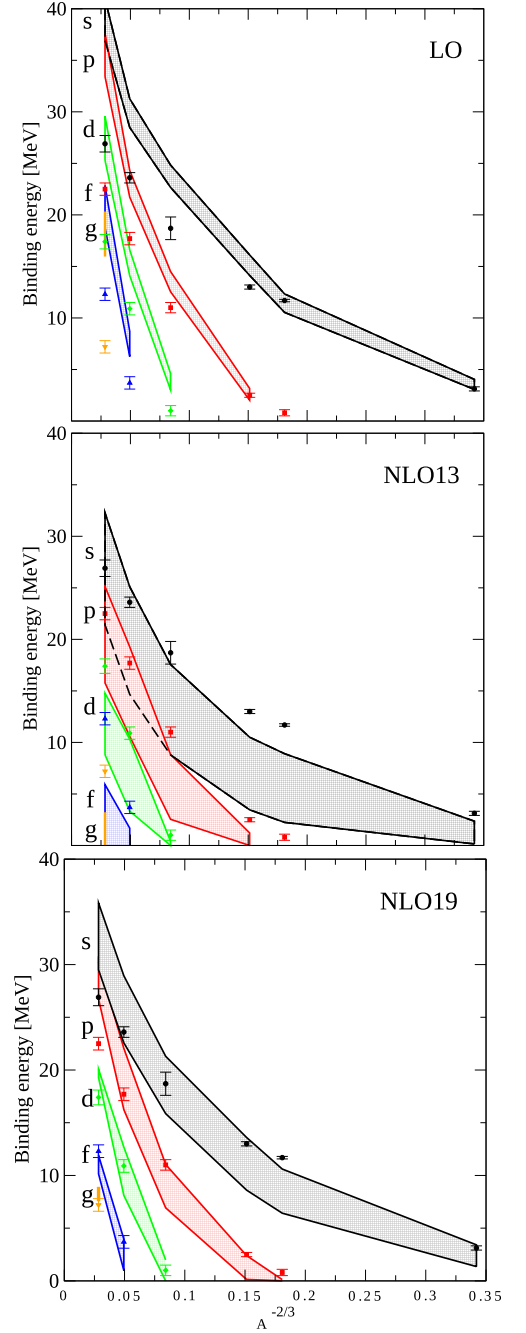


Figure 1.7.: Calculation for the 1s, 1p, 1d, 1f, and 1g Λ single particle shells as a function of $A^{-2/3}$ where A is the mass number of the hypernucleus. The symbols represent experimental results. The figure is reproduced with permission from Springer Nature [28].

calculations.

Furthermore, theoretical calculations within the framework of ChEFT for hypernuclear systems are computationally challenging due to the intricate and unresolved nature of YN interactions, coupled with the complexity of the many-body Hamiltonian. Ab initio methods, such as the No Core Shell Model (NCSM), offer a viable approach to address these challenges. The NCSM simplifies calculations by employing harmonic oscillator (HO) basis states and truncating the Hilbert space to a manageable number of baryonic states [35]. While the model has been successfully adapted to study hypernuclei as exemplified in Figure 1.5, its applicability is currently limited to primarily p-shell hypernuclei and restricted to light hypernuclei, including $^{\Lambda}\text{He}$ and $^{\Lambda}\text{Li}$ isotopes [24]. The extension of the NCSM towards medium-mass hypernuclei, such as $^{20}_{\Lambda}\text{O}$ [36] and $^{41}_{\Lambda}\text{Ca}$ [37], is still in progress.

1.1.3. Nuclear Equation of State and Hyperons

The equation of state (EoS) of the nuclear matter describes the relation between energy (E), density (ρ), pressure (P), temperature (T) and the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$ where ρ_n and ρ_p are neutron and proton densities. At zero temperature, the energy per nucleon is given by

$$E(\rho, \delta) = E_0(\rho, 0) + E_{sym}(\rho)\delta^2 + \mathcal{O}(\delta^4), \quad (1.12)$$

where E_0 is the energy per nucleon of the symmetric nuclear matter (SNM) when $\rho_n = \rho_p$ and E_{sym} is the symmetry energy. Around the nuclear saturation density $\rho_0 = 0.16 \text{ fm}^{-3}$ ($2.8 \times 10^{14} \text{ g/cm}^{-3}$), the first order term of Equation 1.12 can be expanded as

$$E_0(\rho) = E_0(\rho_0) + \frac{K_0}{2!}\chi^2 + \frac{Q_0}{3!}\chi^3 + \mathcal{O}(\chi^4), \quad (1.13)$$

where $\chi = (\rho - \rho_0)/3\rho_0$. The first term $E_0(\rho_0)$ on the right side of Equation 1.13 is the maximum binding energy per nucleon in SNM, which empirically equals approximately 16 MeV. The expansion coefficients, namely the incompressibility coefficient K_0 and the skewness coefficient Q_0 of SNM in higher order terms, are isoscalar parameters that stay invariant under isospin transformation. They are given by

$$K_0 = 9\rho_0^2 \left. \frac{d^2 E_0(\rho)}{d\rho^2} \right|_{\rho=\rho_0}, \quad (1.14)$$

$$Q_0 = 27\rho_0^3 \left. \frac{d^3 E_0(\rho)}{d\rho^3} \right|_{\rho=\rho_0}. \quad (1.15)$$

The incompressibility coefficient is determined experimentally from measurements of nuclear giant monopole resonances. Reported values include $K_0 = 210 \pm 30 \text{ MeV}$ from Ref. [38], $K_0 = 230 \pm 40 \text{ MeV}$ given by Ref. [39] and $K_0 = 230 \pm 10 \text{ MeV}$ suggested by Ref. [40]. The skewness coefficient is often constrained by the flow data analysis of heavy ion collisions (introduction of heavy ion collisions is in Section 1.2.3) and observations of neutron star mass within a range of $-464 \text{ MeV} \leq Q_0 \leq -342 \text{ MeV}$ as detailed in Ref. [41].

While the EOS of SNM is characterized by isoscalar parameters, namely E_0 , K_0 and Q_0 in Equation 1.13, the EOS of asymmetric nuclear matter or pure neutron matter (PNM), which is an approximation of the neutron star matter, is described by a group of isovector parameters (variant under isospin transformation), namely S_0 , L and K_{sym} , as given within the Taylor expansion:

$$E_{sym}(\rho) = S_0(\rho) + L(\rho)\chi + \frac{K_{sym}}{2!}\chi^2 + \mathcal{O}(\chi^3), \quad (1.16)$$

where S_0 is the symmetry energy coefficient, L represents the slope parameter of the symmetry energy and K_{sym} represents the symmetry energy curvature parameter. In contrast to isoscalar coefficients characterizing symmetric nuclear matter, these isovector parameters remain less constrained. Nonetheless, experimental efforts to determine these quantities have been undertaken, such as neutron skin thickness measurements [42], isobaric-analog-state (IAS) phenomenological study incorporating neutron skin data [43], and gravitational observations of neutron stars [44]. For instance, it was proposed in Ref. [43] $S_0 = 31.7 \pm 3.2$ MeV at the saturation density and in Ref. [45] based on neutron star observations, $L = 57.7 \pm 19$ MeV and $K_{sym} = -107 \pm 88$ MeV were suggested.

Precise determination of these parameters is indispensable for understanding the nuclear matter EOS and constraining theoretical models. A diverse range of theoretical frameworks are proposed for computing the EOS, for instance, the Skyrme-Hartree-Fock (SHF) approach [46]. A phenomenological effective nucleon-nucleon (NN) interaction, the Skyrme interaction, characterized by two- and three-body terms, is adjusted to reproduce experimental observables such as nuclear binding energies. With this effective interaction, nucleons are approximated as single-particle wave functions evolving in a self-consistent mean field determined by the Hartree-Fock equations. Alternatively, the relativistic mean field (RMF) theory [47] employs a Lagrangian density to describe nucleons-meson coupling, incorporating nonlinear meson self-interactions as well. Figure 1.8 presents EOS calculations for both symmetric nuclear matter (SNM) and pure neutron matter (PNM) using both SHF and RMF models. While both models accurately reproduce the minimum binding energy of the SNM at saturation density, discrepancies emerge at higher densities, highlighting the need for additional constraints obtained from terrestrial experiments, including heavy-ion collisions, and astronomical observations to refine the comprehension of the EOS in the high-density regime.

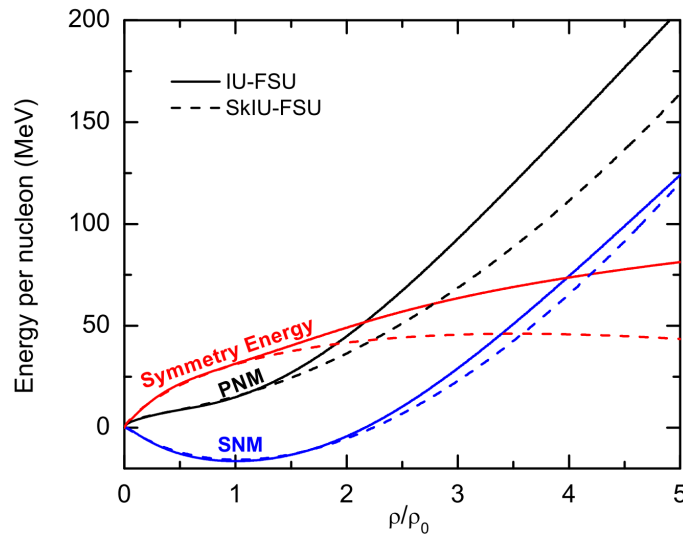


Figure 1.8.: The EOS of SNM (blue) and PNM (black) as well as the symmetry energy (red) as a function of density obtained within the IU-FSU RMF model and the SHF approach using the SkIU-FSU parameter set. Reprinted figure with permission from [48] Copyright 2013 by the American Physical Society.

The inclusion of hyperons in nuclear matter significantly complicates the calculation and modeling of the EOS. Hyperons substantially increase the total density, often surpassing the saturation density. The incorporation of additional two-body (YN and YY) and three-body (YNN) interactions, not readily modeled as nucleon-nucleon interactions, is required. Moreover, the non-zero isospin of Σ and Ξ hyperons introduces more consideration

of the symmetry energy term. Their abundances depend on E_{sym} and fractions of the different hyperonic species, in turn, reflect the symmetry energy behavior [49].

The presence of hyperons in nuclear matter also reduces the stiffness of the EOS, where stiffness characterizes the resistance of nuclear matter to external compression. The additional degrees of freedom associated with hyperons result in a lower energy per baryon when the baryon density increases. Moreover, hyperons can occupy low-energy states, thereby reducing the Fermi pressure exerted on baryons. Figure 1.9(a) displays the effect of the presence of hyperons on the pressure where the energy density ϵ in a dense environment, such as in the core of a neutron star, can be approximated to [50]

$$\epsilon = \rho c^2. \quad (1.17)$$

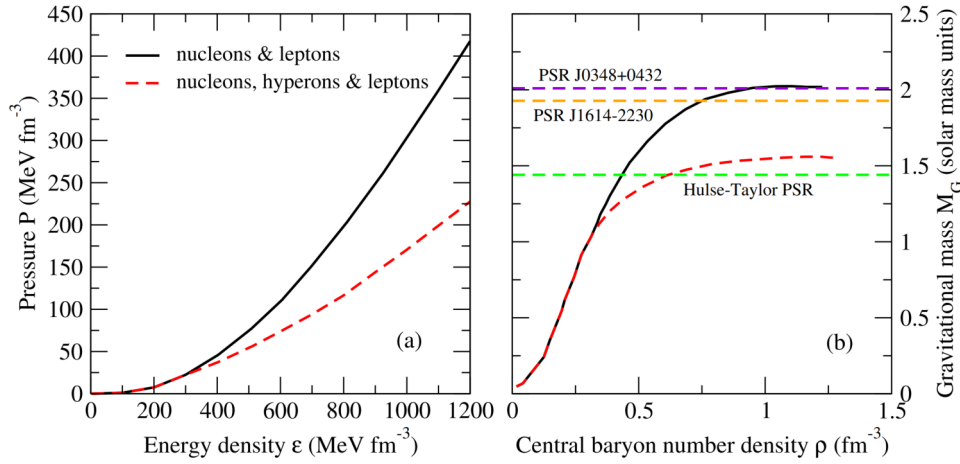


Figure 1.9.: (a) The EOS pressure as a function of energy density with and without the presence of hyperons and (b) masses of neutron stars. A generic model with (black solid line) and without (red dashed line) hyperons has been considered. The horizontal lines show the observational mass of the Hulse–Taylor pulsar [51] and the recently observed PSR J1614-2230 [52] and PSR J0348+0432 [53]. Used with permission of The Royal Society (U.K.), from [54]; permission conveyed through Copyright Clearance Center, Inc.

This softening effect contributes to the so-called ‘Hyperon Puzzle’. In the core of a neutron star, the density reaches 2-3 ρ_0 [55] and even beyond that. The pressure of the neutron star matter in this region is the Fermi pressure which counterbalances the inward gravitational pressure, preventing the entire neutron star from collapsing. However, at such high densities, the nucleon chemical potential is sufficient to convert nucleons into hyperons [56], first Σ^- followed by Λ [57], which is energetically unavoidable. This leads to a decreased pressure, i.e., the softening of EoS for neutron star matter. The softening effect subsequently results in a reduction of the maximum mass of the neutron star. The relationship between the pressure and the total mass of the neutron star is given by the Tolman-Oppenheimer-Volkof (TOV) equation [58],

$$\begin{aligned} \frac{dM(r)}{dr} &= 4\pi\epsilon(r)r^2, \\ \frac{dP(r)}{dr} &= -\frac{[\epsilon(r) + P(r)][M(r) + 4\pi P(r)r^3]}{r^2[1 - 2M(r)/r]}, \end{aligned} \quad (1.18)$$

where M represents the mass and ϵ represents the energy density. The units are adopted with gravitational constant $G = 1$ and the speed of light $c = 1$. Neutron star masses can be calculated from the Equation 1.18

at a specified energy density. The predicted mass estimate can subsequently be compared with observational data for validation.

It is illustrated in Figure 1.9(b) that the softening of the EoS is incompatible with the observation of massive pulsar PSR J1614-2230 ($1.928 \pm 0.017 M_{\odot}$) [52] and PSR J0348+0432 ($2.01 \pm 0.04 M_{\odot}$) [53]. The calculated mass with the presence of hyperons is below the observed mass of both neutron stars.

To reconcile the contradiction between the observed masses of neutron stars and the softening of the EOS by hyperons, additional repulsion is required to stiffen the EOS of the neutron star matter. One possible mechanism within the hypernuclear framework is the incorporation of repulsive hyperon-hyperon (YY) interactions, which could potentially support neutron star masses up to $2 M_{\odot}$ [59]. Alternatively, hyperon-nucleon three-body interactions (YNN, YYN, and YYY) might provide the necessary stiffness [60]. While the precise nature of these interactions remains to be determined through both experimental and theoretical efforts, the study of hypernuclei offers a valuable approach to probing them.

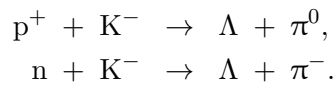
1.2. Production of Λ -Hypernuclei in Experiments

Hypernuclei serve as an important probe to understand YN and YY interactions, while experiments with hypernuclei are technically challenging compared to normal nuclei. On one hand, direct scattering and reaction experiments are extremely difficult owing to the short lifetime of hypernuclei since it is impossible to produce a stable and intense hypernuclei beam. In most facilities, hypernuclei are produced in a stable target instead. On the other hand, high-precision detection techniques are necessary to track down and reconstruct the decay of the hypernuclei.

Owing to the relatively longer lifetime of Λ than other hyperons, Λ hypernuclei have been the most investigated systems and the cornerstone of strangeness nuclear physics. To produce Λ hypernuclei, there are in general several directions: (1) strangeness exchange reaction, (2) strangeness production reaction, and (3) heavy ion collisions.

1.2.1. Strangeness Exchange Reaction

It is a natural method to impinge on a nuclear target with a beam carrying strangeness directly, in most cases, kaon beams. a K^- ($s\bar{u}$) meson collides with a nucleon in the target and undergoes the strangeness exchange reaction where the up or down quark in the nucleon is replaced by the strange quark in the K^- , as shown in the following:



Both reactions are exothermic with a Q -value of approximately 181 MeV and 178 MeV, respectively. The reaction cross-section ratio of (K^-, π^0) to (K^-, π^-) is 0.5 due to the $\Delta I = 1/2$ rule [61] which summarizes that experimentally the branching ratio of charged channels is as twice large as that of neutral channels. This rule holds for other non-leptonic strangeness-changing processes [62], like the decay of the Σ hyperon and pionic Kaon decays [63]. Depending on the energy of incoming K^- , the reaction mechanism differs. For (K_{stop}^-, π) reaction, Kaons are absorbed from the atomic orbit into the nucleus accompanying the emission of X-rays. The production rate of the Σ is higher than the Λ in this scenario. The momentum transferred to the recoil hyperon is close to the Fermi momentum, thus this method exhibits high quasi-free Σ background [64]. In addition, the energy resolution is limited since a thick target slows down the beam to generate stopped Kaons [65].

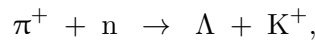
The in-flight $(K_{\text{in-flight}}^-, \pi)$ reaction presents lower momentum transfer. When the incident beam momentum is

around 505 MeV/c, the momentum transfer at 0° is considerably smaller than the Fermi momentum k_f even to zero [66], thus the Λ hyperon converted from the nucleon remains in the same orbit. One advantage is that by selecting the beam momentum to 700 - 800 MeV/c, the reaction cross-section reaches its maximum [67]. The strangeness exchange reaction is massively studied by BNL [68][69] and KEK [70].

1.2.2. Strangeness Production Reaction

Pion-Induced Associated Production

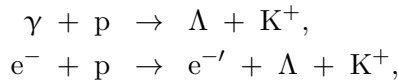
The associated production reaction employs the π^+ beam. The nucleon interacts with the incoming pion leading to an excitation of an intermediate baryon resonance, either Δ^* , or N^* , with a typical mass of the Δ^* or N^* resonance of 2 GeV/c² [71]. Subsequently, the resonances, particularly $N^*(1650)[\frac{1}{2}^-]$, $N^*(1710)[\frac{1}{2}^+]$ and $N^*(1720)[\frac{3}{2}^+]$ [72], decay into ΛK^+ . The overall reaction is given by



with a Q -value of approximately -530 MeV. The momentum and angular momentum transferred to Λ are higher than in the strangeness exchange reaction. When at the beam momentum of 1.04 GeV/c, the outgoing K^+ has a momentum of approximately 0.7 GeV/c while Λ has a momentum of approximately 300 MeV/c for light hypernuclei [73] which is larger than the Fermi momentum. Therefore the Λ can populate a wide range of orbits without taking up a nucleon hole. Consequently, it reveals the shell structure of Λ hypernuclei [74]. Another big advantage of this reaction is that it permits the excitation of high-spin hypernuclear states. However the cross-section of the associated production reaction is smaller than strangeness exchange reactions [65], yet high-intensity pion beams are easier to obtain than Kaon beams. This reaction is studied by FINUDA [75], KEK[76], and J-PARC E40[77].

Electromagnetic Production

An alternative method that offers better energy (several hundred keV) and spatial resolution than hadron-induced reactions is the photoproduction or electron-induced production [78], which is also associated production, with the following process



with a Q -value of approximately -671 MeV. A real photon or a virtual photon created by an electron excites the proton to a Δ^* which then decays into ΛK^+ .

The hyperons produced in electromagnetic production have high momentum and angular momentum transfer indicating a high-spin hypernucleus state [79]. The momentum transfer is greater than 350 MeV/c thus there is a probability of hyperons escaping from the nucleus, while the cross-section is lower than pion-induced associated production [65] though a higher intensity electron beam can compensate for that.

There are several other advantages of electromagnetic production. Photons interact weakly with the nucleus thus allowing the production of deeply bound hypernuclei. The transition of this associated production can be employed to probe spin-flip states [80]. Furthermore, the reaction transforms the proton to a hyperon thus allowing the production of pairs of neutron-rich mirror hypernuclei by using the same target [81].

The electromagnetic production method is employed in experiments at Jlab [82][83] and Mainz [84].

1.2.3. Heavy Ion Collisions

In heavy ion collision (HIC) experiments, nucleons of heavy nuclei, for example, lead (Pb) and gold (Au), collide at relativistic energies to search for new forms of matter in extreme conditions and to explore the

possibility of producing Quark-Gluon Plasma (QGP) which is an extreme state of matter with ultra-high temperature and energy density where color charges are not confined within hadrons. QGP is believed to have existed for a few millionth of a second after the Big Bang.

General Steps of Heavy Ion Collisions

The temporal evolution of an event in HIC experiments can be described by the following stages (Figure 1.10).

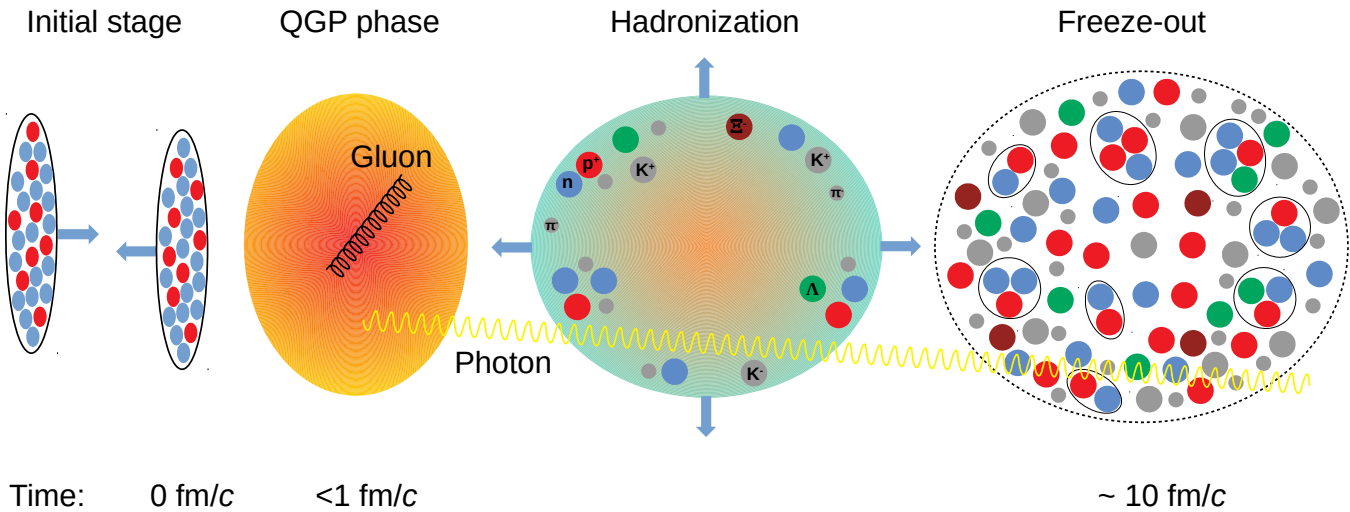


Figure 1.10.: Different stages of an event from HIC experiments.

1. Initial state: Two nuclei approach at a center-of-mass energy, so-called collision energy $\sqrt{s_{NN}}$, in a way that can be visualized as Lorentz-contracted discs of nucleons. These discs, due to relativistic effects, appear flattened in the direction of motion. The contraction for heavy nuclei is of the order of $14/\gamma$ fm [85] (γ is the Lorentz factor). As they collide, the nucleons within these discs interact non-uniformly, leading to a complex and dynamic collision environment.
2. QGP phase: Two nuclei overlap and produce a volume of high temperature and energy density which is known as the fireball. Inside the region, the matter may undergo a phase transition to QGP. A process called thermalization [86] is reached provided that the density and temperature inside the fireball are substantially high for sufficient strong interactions to take place, resulting in a thermodynamic equilibrium of quarks and gluons. The sufficiently high energy density and the temperature allow these quarks and gluons to overcome the binding strong force leading to the deconfinement.
3. Hadronization: The fireball expands due to the high pressure in the center resulting in a decrease in the density and the temperature. When the temperature drops to a certain level during the cooling of the QGP, quarks and gluons start to recombine and form particles like hadrons. During the expansion phase, a collective phenomenon of particle emission, known as the flow, is characterized by an anisotropic distribution. It arises due to initial pressure gradients and subsequent particle interactions [87]. The evolution of the flow is found to be sensitive to the stiffness of the EOS [88] of quark and nuclear matter.
4. Freeze-out: The expansion continues until particle inelastic interactions cease and the hadron abundances stop changing. This point is called chemical freeze-out. At a later point, all types of momentum-changing elastic and inelastic interactions cease, which is called kinetic freeze-out. Depending on the particle

species, the density, and flow velocity, different parts of the fireball decouple at different times [89]. The lifetime of the fireball, which spans from the pre-collision stage to the freeze-out stage, varies on the order of 10 fm/c (approximately 3×10^{-23} s) depending on $\sqrt{s_{NN}}$ [90].

For the aforementioned reaction mechanisms, the production of hypernuclei is confined in the stability valley of nuclei since these reactions only convert one nucleon of the nucleus in the stable target to a hyperon. HIC experiments on the other hand enable the production of hypernuclei far from the stability line [91].

Strangeness Production

The dominant mechanisms for strangeness production in HICs depend on $\sqrt{s_{NN}}$. At the high energies of experiments like the Large Hadron Collider (LHC) [92] and the Relativistic Heavy Ion Collider (RHIC) [93], the formation of the QGP is widely accepted. Within this state, the production of strange quark-antiquark pairs occurs substantially, followed by the formation of strange mesons and hyperons during the hadronization process. Thus strangeness enhancement is believed to be one crucial evidence for QGP formation [94]. More details related to these experiments can be found in [95].

However, at the current FAIR SIS18 energy regime, typically below 2 AGeV[96], the possibility of QGP formation is less certain [97]. Quark confinement is probably not overcome at these energies, thereby prohibiting the production of strange quark-antiquark pairs. Instead, associated production in nucleon-nucleon interactions remains the primary mechanism for strangeness production. The incident energy threshold of NN elementary collisions for strangeness production[98] is 1.58 AGeV:



It is also probable to produce ΛK at sub-threshold energy either through multiple NN collisions creating intermediate $N\Delta$ or $\Delta\Delta$ pairs and subsequently ΔN (or $\Delta\Delta$) $\rightarrow N\Lambda K^+$, or via secondary $\pi^+ n \rightarrow \Lambda K^+$, where the π^+ is produced from NN interaction. The down quark from the neutron and the $\bar{d}$ quark annihilate to form a $s\bar{s}$ pair. Both mechanisms are significantly influenced by in-medium effects [99] since with a high baryon density, the probability of a cascade of NN collisions is high, accumulating the energy and catalyzing the production of ΛK at a sub-threshold energy.

Hypernuclei Formation

The predominant production of light clusters and fragments in the freeze-out stage is observed from emitted particles in HIC experiments [100]. However, the formation mechanism of hypernuclei remains elusive due to their scarcity in experiments, leading to deviation between different theoretical models.

Coalescence models with the hadronic transport model [101] propose that hypernuclei production occurs via the absorption of hyperons by spectator fragments in peripheral collisions [102]. The hypernuclear formation in the proximity of nucleons and hyperons in momentum and coordinate space is postulated, while due to strangeness suppression [103] hypernuclei production is characterized by low yields. Consequently, the coalescence parameters are not well constrained and the model is thus sensitive to the binding energy of hyperons [104].

In contrast, thermal models assume that during the expansion stage, the hot nuclear matter undergoes a liquid-gas phase transition [105]. This transition is considered to facilitate the formation of light clusters, characterized by spatial correlations between multiple nucleons and a Λ hyperon. Without delving into the internal structure of hypernuclei, statistical approaches are employed in thermal models to calculate the expected yield of those hypernuclei in a thermal equilibrium conserving the entropy per baryon [106].

To compare these two models, observables, such as hypernuclear yields measured in HIC experiments, offer a valuable means of assessment, although they are subject to experimental acceptance. To mitigate the impact of acceptance and efficiency corrections, the yield ratio of a hyperon to its corresponding hypernucleus, or that of a hypernucleus to a nucleus, such as the ratio of ${}^3_\Lambda\text{H}/{}^3\text{He}$, is always employed. The coalescence model

reproduced the yield ratio ${}^3_\Lambda\text{H}/{}^3\text{He}$ from STAR experimental results obtained from gold-gold collisions at a collision energy of 200 GeV with a centrality range of 0-80%, while the thermal model exhibits better alignment with the yield ratio from data obtained by the ALICE collaboration collected from lead-lead collisions at 2.76 TeV within a centrality range of 0-10% [107]. The discrepancy between the two models in their respective agreement with STAR and ALICE data may be attributed to a dependence on the initial conditions of the collision, such as $\sqrt{S_{\text{NN}}}$. Both models described the yield of ${}^3_\Lambda\text{H}$ at $\sqrt{S_{\text{NN}}} = 3$ GeV while underestimated the yield of ${}^4_\Lambda\text{H}$ [108]. To better determine the production mechanism within HIC, precise measurements of the lifetime, yield, size, and binding energy of hypernuclei are critical [109].

The investigation of hypernuclei production is extensively conducted in a wide range of HIC experiments, such as STAR [110], ALICE [111], and HyPHI [112]. They provide a great opportunity to produce a wide variety of hypernuclei with different compositions, although hypernuclei production in these HIC experiments suffers from the low yield which introduces uncertainty.

1.3. Hypernuclear Halos and the Hypertriton Case

1.3.1. Halo Nucleus

Definition of a Halo Nucleus

A halo nucleus is a nuclear system with a valence nucleon weakly bound to the core while this nucleon occupies the state shallowly below the surface of the potential well. Due to the tunneling effect, its wave function extends beyond the well walls. Consequently, there is a large probability of finding the nucleon outside the well. By definition, a halo nucleus is required to have more than 50% of its probability density outside the range of the core potential [113]. In the case of a neutron halo, its wave function $\psi(r)$ outside the potential well can be approximated with the Yukawa type,

$$\psi(r) = A \frac{e^{-\kappa r}}{\kappa r}, \quad (1.19)$$

$$\kappa = \frac{\sqrt{2\mu S_n}}{\hbar}, \quad (1.20)$$

where r denotes the distance from the center of the nucleus, A is the normalization constant, and μ represents the reduced mass of the core-neutron system. This Yukawa-type wave function describes the neutron's asymptotically decaying behavior. As the radial distance from the nuclear center increases, the wave function exhibits an exponential decay that is inversely dependent on the neutron separation energy S_n . A lower S_n corresponds to a longer decay tail. This separation energy-dependent decay behavior can be characterized by the mean square radius of the nucleus $\langle r^2 \rangle$, which is defined by

$$\langle r^2 \rangle = \frac{\int r^4 \psi^2(r) dr}{\int r^2 \psi^2(r) dr} = \frac{\hbar^2}{4\mu S_n}. \quad (1.21)$$

It can be noted from Equation 1.21 that the root mean square (rms) radius $\sqrt{\langle r^2 \rangle}$ is inversely proportional to the square root of S_n . A nucleus exhibiting low separation energy, typically less than approximately 1 MeV, is often indicative of a halo nuclear structure and thus shows a larger rms radius. For instance, ${}^{19}\text{C}$, experimentally determined to be a one-neutron halo nucleus, has a lower neutron separation energy $S_n = 242 \pm 95$ keV and a larger estimated rms radius of 6.0 ± 0.9 fm [114] compared to ${}^{17}\text{C}$, which has a higher neutron separation energy $S_n = 734 \pm 18$ keV and a smaller radius of 3.0 ± 0.6 fm [115].

Furthermore, the rms radius depends on the angular momentum (l) of the orbital. It is suggested that the

halo phenomenon is associated with nuclei with low orbital angular momentum ($l \leq 1$) states since at low separation energy ($S_n \rightarrow 0$), the rms radius diverges only for s- ($l = 0$) and p- ($l = 1$) wave states [116], as illustrated in Figure 1.11.

Radius Determination of a Halo Nucleus

The radius of a halo nucleus can be determined experimentally. One effective approach involves the application of the Glauber theory. This theoretical framework simplifies the modeling of nucleus-nucleus collisions by employing the semi-classical eikonal approximation [117].

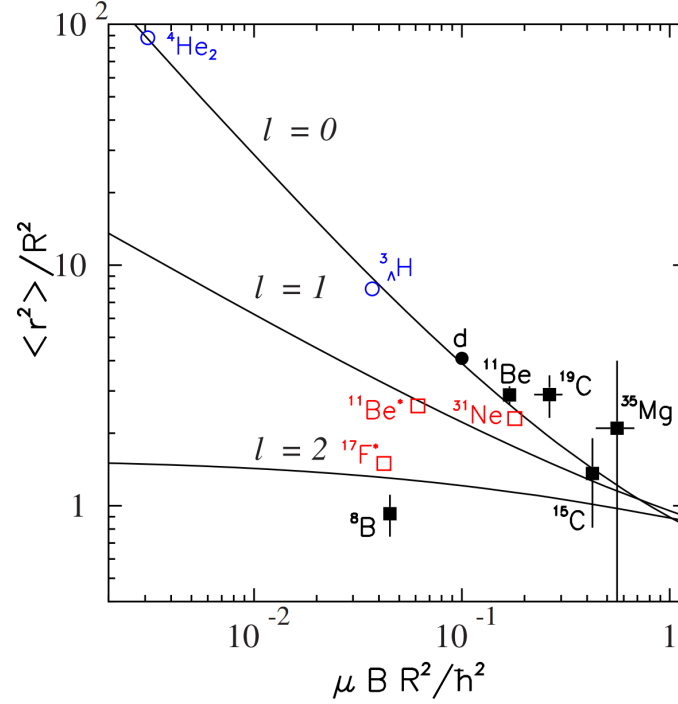


Figure 1.11.: Dimensionless mean square of radii and dimensionless two-body binding energy for halo systems. R represents the classical turning point and B represents the separation energy of the valence nucleon or hyperon. The filled circle denotes the deuteron, the filled squares denote nuclei where radii were extracted from experimental interaction cross-sections, the open squares are simple model estimates and the open circles are theoretical calculations. Used with permission of IOP Publishing, Ltd, from [118]; permission conveyed through Copyright Clearance Center, Inc.

This approximation postulates the following conditions with the incident wave number, k , the range of the interaction, a , the incident particle's energy, E , and the potential strength, V_0 :

$$\begin{aligned} ka &\gg 1, \\ E &\gg |V_0|. \end{aligned} \quad (1.22)$$

The approximation requires a considerably smaller de Broglie wavelength of the incident particle than the size of the interaction region and a high scattering energy. In such a case, the nuclei are simplified as disks of nucleons that only travel along definite straight lines. Consequently, the high energy collision between a target nucleus, A , and a projectile nucleus, B , is reduced to a series of elementary nucleon-nucleon interactions as depicted in Figure 1.12. The applicability of the Glauber theory is generally limited to the high-energy

regime according to Equation 1.22, since at lower energies, the Coulomb repulsion can deflect nucleons from straight-line trajectories. The eikonal phase shift function [119] serves as a crucial determinant of nucleon-nucleon scattering cross-sections. It is defined by [120]

$$\chi_{NN}(b) = -\frac{1}{\hbar} \int_{-\infty}^{\infty} V(r, b) dz, \quad (1.23)$$

where $V(r, b)$ is complex and describes the nucleon-nucleon interaction potential with more details in [121][122]. Consequently, the total reaction cross-section can be calculated with

$$\sigma_R = \int (1 - |e^{i\chi_{NN}(b)}|^2) db. \quad (1.24)$$

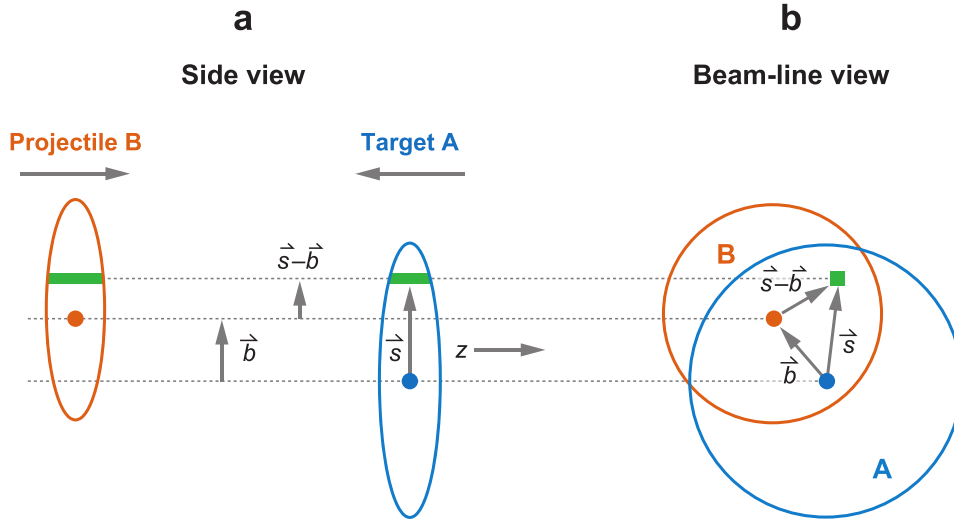


Figure 1.12.: Schematic drawing of the nucleus-nucleus collision in the framework of the Glauber model. The mass numbers of the projectile and the target are A and B . $\vec{b}$ is the vector for the impact parameter. $\vec{s}$ represents the distance between the center of the nucleus A and the interacting flux of nucleons. Used with permission of Annual Reviews, Inc., from [123]; permission conveyed through Copyright Clearance Center, Inc.

For composite nucleus-nucleus collisions, the Glauber theory assumes that the total phase shift function is the sum of the phase shift function of each nucleon-nucleon interaction [124], which is given by

$$\chi(\vec{b}, \vec{s}_1, \vec{s}_2, \dots, \vec{s}_A) = \sum_{j=1}^A \chi_j(\vec{b} - \vec{s}_j), \quad (1.25)$$

where $(\vec{s}_1, \vec{s}_2, \dots, \vec{s}_A)$ are the positions of nucleons of the nucleus A . Therefore the overall phase shift function

for nucleus-nucleus A - B collision can be calculated with [125]

$$\begin{aligned}
 e^{i\chi(\vec{b})} = & \int \cdots \int \left(\prod_{j=1}^A d\vec{r}_j \right) \left(\prod_{k=1}^B d\vec{r}_k \right) \\
 & \times \rho_A(\vec{r}_1, \vec{r}_2, \dots, \vec{r}_A) \rho_B(\vec{r}_1, \vec{r}_2, \dots, \vec{r}_B) \\
 & \times \prod_{j=1}^A \prod_{k=1}^B [1 - \Gamma(\vec{b} + \vec{s}_j^A - \vec{s}_k^B)],
 \end{aligned} \tag{1.26}$$

where $\Gamma(\vec{b} + \vec{s}_j^A - \vec{s}_k^B) = 1 - e^{i\chi(\vec{b} + \vec{s}_j^A - \vec{s}_k^B)}$ is the profile function. The nucleon density function $\rho_A(\vec{r}_1, \vec{r}_2, \dots, \vec{r}_A)$ represents the probability density of finding A nucleons simultaneously at positions $(\vec{r}_1, \vec{r}_2, \dots, \vec{r}_A)$, respectively. Analogously, the nucleon density function $\rho_B(\vec{r}_1, \vec{r}_2, \dots, \vec{r}_B)$ describes the corresponding probability distribution for Nucleus B .

Equation 1.26 demonstrates the direct influence of the nucleon density on the total phase shift function, which subsequently determines elastic scattering and total reaction cross-sections. Monte Carlo simulations are often employed to evaluate the phase shift function and the cross-sections by sampling from the nucleon density distribution. For instance, Figure 1.13 illustrates the calculations for the relationship between the total reaction cross-section of ^{22}C - ^{12}C collisions and the rms radius of ^{22}C , using the Monte Carlo Glauber theory. The accuracy of these calculations depends on the granularity of the nucleon density sampling, i.e., the number of configurations M .

By comparing the calculated reaction cross-section obtained from Monte Carlo Glauber simulations with experimentally measured total reaction cross-section for collisions between a nucleus of interest and a target nucleus with a known nucleon density distribution, the parameters characterizing these density distributions can be constrained. Subsequently, the nucleon density distribution derived from these models is employed to determine the rms radius of the nucleus of interest using Equation 1.21.

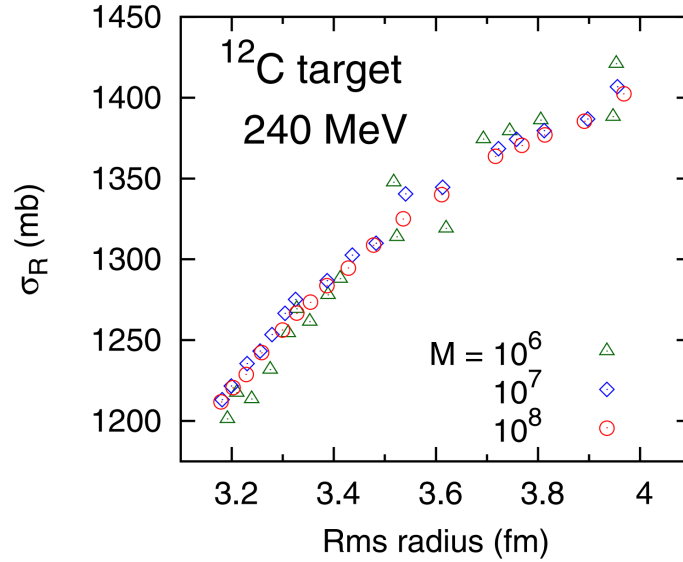


Figure 1.13.: Total reaction cross-section of ^{22}C - ^{12}C collisions as a function of the rms radius of ^{22}C calculated by Monte Carlo Glauber theory. M represents the number of configurations. Reprinted figure with permission from [125] Copyright 2018 by the American Physical Society.

1.3.2. Hypertriton

The hypertriton, denoted as ${}^3_\Lambda\text{H}$, is the lightest hypernucleus, comprising a proton, a neutron, and a Λ hyperon, thus it possesses a spin-parity of $1/2^+$ and an isospin of 0. Due to its uncomplicated composition, the hypertriton serves as a pivotal probe of the YN interaction and hypernuclear structure and is the benchmark for hypernuclear physics.

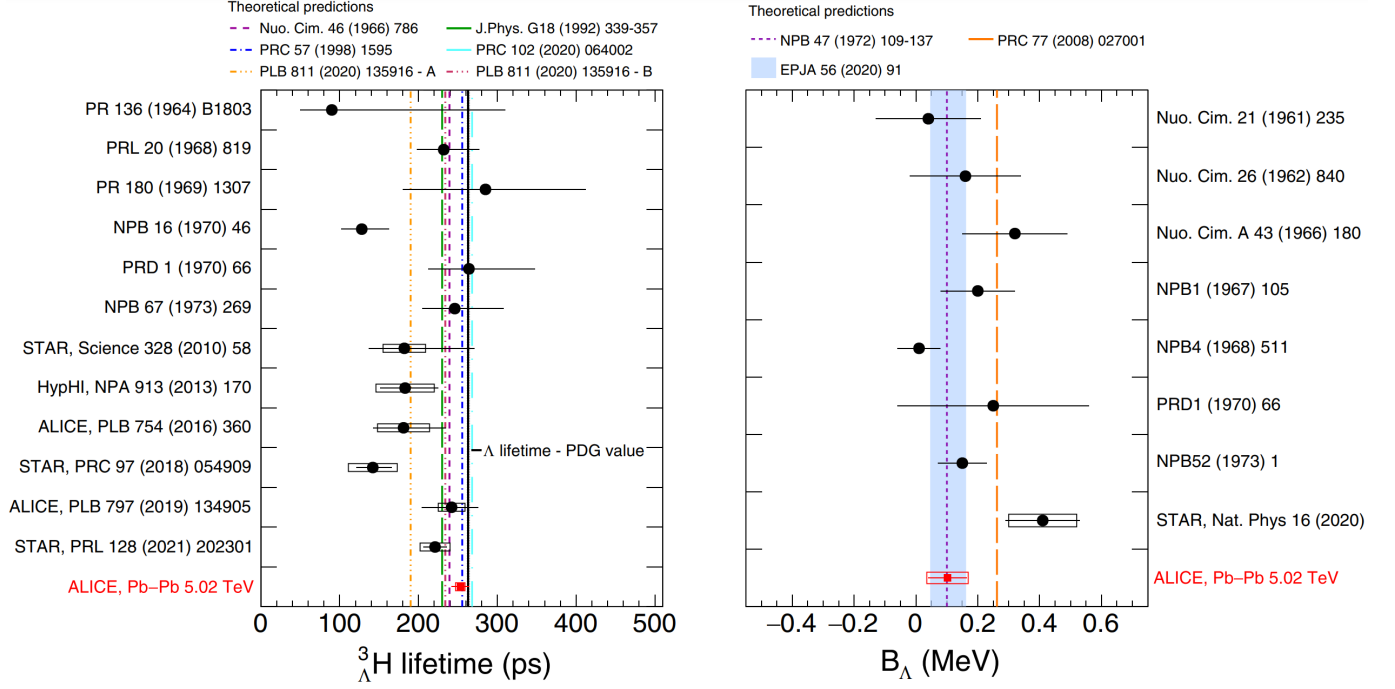


Figure 1.14.: ${}^3_\Lambda\text{H}$ (left) lifetime, τ , and (right) binding energy, B_Λ , measurements obtained with different experimental techniques. The horizontal lines and boxes are the statistical and systematic uncertainties, respectively. The dashed and dash-dotted lines are the corresponding theoretical predictions. The figure is reprinted from [126] under CC BY 4.0.

The hypertriton mainly decays through pionic decay mode, with the following major decay channels [127]:

$$\begin{aligned}
 {}^3\text{He} + \pi^- & \text{ and } {}^3\text{H} + \pi^0 & \Gamma/\Gamma_\Lambda &= 0.384 \\
 \text{d} + \text{p} + \pi^- & \text{ and } \text{d} + \text{n} + \pi^0 & \Gamma/\Gamma_\Lambda &= 0.619 \\
 \text{p} + \text{p} + \text{n} + \pi^- & \text{ and } \text{p} + \text{n} + \text{n} + \pi^0 & \Gamma/\Gamma_\Lambda &= 0.0097 \\
 & & \Gamma_\pi/\Gamma_\Lambda &= 1.01
 \end{aligned}$$

where π^- channels and π^0 channels follow (though slightly violate) the $\Delta I = 1/2$ rule. The total decay rate of NMWD ($\text{d} + \text{p}$ and $\text{p} + \text{p} + \text{n}$) for the hypertriton is found to be 0.5% of the free Λ decay rate [128]. Furthermore, it is suggested to be a weakly bound system, featuring a valence Λ and a deuteron core [129]. Due to these factors, the hypertriton is predicted to exhibit a lifetime comparable to the Λ hyperon, as indicated by pionless EFT [130] calculations, with a binding energy below 0.5 MeV. Experimental efforts have been made to measure both the lifetime and the binding energy, B_Λ of the hypertriton, as presented in Figure 1.14 with a collection of past measurements of these properties. The ALICE collaboration provided high-precision results recently, reporting a hypertriton lifetime of $253 \pm 11(\text{statistical}) \pm 6(\text{systematic})$ ps [126], which is in excellent agreement with theoretical predictions. Additionally, they determined a binding

energy, B_Λ , of $102 \pm 63(\text{statistical}) \pm 67(\text{systematic})$ keV for the hypertriton [126]. While the central value of this measurement aligns with some theoretical predictions, which evidences the weakly bound feature of the hypertriton, it is subject to significant uncertainty.

Due to its weakly bound nature, the hypertriton is predicted to be a hyper-halo nucleus. As suggested by Equation 1.21, the binding energy is approximately inversely proportional to the square root of the rms radius of the hypertriton. The matter radius of hypertritons is inferred to be 4.2 fm using ChEFT calculations within the NCSM formalism [131] corresponding to the binding energy of 130 keV. A lower binding energy leads to a maximum radius of 10 fm as shown in Figure 1.11. Thus accurate determination of the rms radius of the hypertriton can contribute to reducing the uncertainty of the binding energy. However, no experimental data for the radius of hypertritons is available yet.

The size of the hypertriton also sheds light on nuclei formation in HIC experiments. While the thermal model does not involve the internal structure of hypertriton, the coalescence model on the other hand takes the nuclear wave function as one input that determines the formation probability of bound states [132]. Since the size of the hypertriton wave function (4-10 fm) is larger than the nucleon and hyperon emission radius (1-2 fm) [133], the yield of hypertritons in HIC experiments predicted by the coalescence model is suppressed [134] since the probability of finding the necessary nucleons and hyperons to interact and form the hypertriton is lower for a larger nucleus size to source size ratio.

1.4. Introduction to HYDRA Experiment

HYDRA (HYpernuclei Decay at R³B Apparatus) is a physical program that aims for hypernuclear invariant-mass spectroscopy and light hypernuclei study via HIC at GSI/FAIR energies. At phase 0, the experiment, with a 1.9 GeV/u ¹²C⁶⁺ beam impinging onto a ¹²C target, is dedicated to the determination of the hypertriton radius. The interaction cross-section between ¹²C and the hypertriton, produced at ~ 1.5 GeV/u in ¹²C-¹²C collisions, is measured.

1.4.1. Challenges and Solutions for Hypertriton Experiments

Experimentally it is challenging to perform direct measurements of the hypertriton's radius or its interaction cross-section with another target. Owing to its sub-nanosecond lifetime and low production cross-sections, as suggested by Ref. [91] that hypertriton production cross-section from ¹²C-¹²C collisions is 1.8 μb . It is infeasible to either produce a stable and sufficient intensity hypertriton beam or to use a fixed hypertriton target. A two-step approach, described in [135] with more details, was developed for studying unstable nuclei. By employing this technique, hypertritons are initially produced from HIC collisions, followed by a second step where the hypertriton can interact with the target before decaying.

Furthermore, owing to its short lifetime, the hypertriton decays very soon after production thus it is not possible to directly detect hypertritons. Instead, its product particles are always measured by detectors. In this case, a so-called invariant mass technique is usually employed to study hypertritons. The invariant mass of the hypertriton $m_{\Lambda\text{H}}^3$ can be computed from

$$m_{\Lambda\text{H}}^2 = (\sum E_i)^2 - \left(\sum \vec{p}_i\right)^2, \quad (1.27)$$

where E_i and $\vec{p}_i$ are the energy and the momentum of the daughter particle i from the hypertriton decay, respectively. The total yield of hypertritons is determined by counting the number of events within a specified invariant mass range centered on the nominal mass of the hypertriton. However, there is an additional source of uncertainty that stems from background contamination. In ¹²C-¹²C collisions, the production of multiple pions or protons within one event may significantly hinder the reconstruction of the hypertriton invariant mass

spectrum. This hindrance arises from the mismatching of non-correlated decay products when the hypertriton decays within a distance from the target that is smaller than the detector's resolution for distinguishing decay products. In such cases, meson beams are not favored due to the fact that the energy of the meson-induced produced hypernuclei is always low, as discussed in Section 1.2.1 and 1.2.2, which is not able to separate the vertex of hypertriton decay from the primary ^{12}C - ^{12}C collision vertex.

To confront the aforementioned challenges, a promising strategy involves dedicated experiments for hypernuclear invariant-mass spectroscopy utilizing ion beams and fixed targets. This approach was first implemented by the HypHI Phase 0 experiment at GSI [136], where a 2 GeV/nucleon ^6Li beam was incident upon a 4 cm ^{12}C target. Owing to the substantial Lorentz factor, the production and decay vertices are spatially separated by 10-30 cm. This effect, referred to as the off-decay vertex, enables the hypernucleus to traverse a measurable distance downstream of the target before decaying.

The HYDRA experiment adopts the same strategy with the beam energy of 1.9 GeV/u which exceeds the incident threshold energy of strangeness production. Considering the hypertriton's lifetime of 253 ps as reported by the latest ALICE results, hypertritons can travel on average 7.2 cm following its production before decaying. With the help of the high-acceptance magnet GLAD (detailed in Chapter 3) operating at ~ 2 T, a momentum resolution of 1% for the decay daughter particles is expected, leading to a decay vertex resolution better than 5 mm and a mass resolution of 1.5 MeV for the hypertriton [137].

1.4.2. Experimental Method of HYDRA

When the hypertriton is produced from ^{12}C - ^{12}C collision in the target, there are mainly two follow-up processes: (1) the hypertriton decays through the channels described in Section 1.3.2 and (2) the hypertriton interacts with target particles. The production cross-section (σ_Λ) of hypertritons is unknown as well as the interaction cross-section ($\sigma_{\Lambda R}$) of hypertritons and target particles. Two independent measurements with different target thickness are needed to calculate $\sigma_{\Lambda R}$. The "Two-Target" Method is thus proposed to extract the value.

Assuming propagation of particles along the beam axis (z), the population $N_\Lambda(z)$ of any hypernuclei $^A_\Lambda\text{X}$ can be given by [138]

$$dN_\Lambda(z) = -\frac{N_\Lambda(z)}{\gamma\beta c\tau}dz + \delta(n\sigma_\Lambda N(z) - n\sigma_{\Lambda R}N_\Lambda(z))dz, \quad N_\Lambda(0) = 0, \quad (1.28)$$

$$dN(z) = -\delta(n\sigma_R N(z))dz, \quad N(0) = It, \quad (1.29)$$

where $\delta = 0$ or 1 (0 in free space and 1 in the target). β is the velocity of $^A_\Lambda\text{X}$ in the laboratory frame, γ its Lorentz factor and τ is $^A_\Lambda\text{X}$ lifetime. n is the target particle number density while $N(z)$ is the population of the beam particles. I is the beam intensity and t is the total period. Equation 1.28 consists of three parts, weak decay of $^A_\Lambda\text{X}$, production of $^A_\Lambda\text{X}$, and interaction of $^A_\Lambda\text{X}$ with the target particles, respectively.

The sensitivity of the $^A_\Lambda\text{X}$ yield to $\sigma_{\Lambda R}$ is displayed in Figure 1.15. In the case of the hypertriton, its decay daughter particles from the two-body decay channel, namely the π^- and the ^3He , are tracked inside the magnet by HYDRA-TPC and fiber detectors, respectively with more details in Section 3.1.1. Following track reconstruction, the radius of curvature of each particle's trajectory is extracted. In conjunction with the known magnetic field strength, the magnetic rigidity, $B\rho$, of the particle is determined. Subsequently, the momentum p of the particle can be calculated using the relation:

$$p = qB\rho, \quad (1.30)$$

where q is the particle charge, B is the magnetic flux density and ρ is the density. The invariant-mass spectrum of the hypertriton from the two-body decay channel can therefore be reconstructed using Equation 1.27.

Consequently, the total yield of hypertritons at the target exit point, $N_{\Lambda}(d)$, can be inferred from the two-body decay yield and its branching ratio. The interaction cross-section, $\sigma_{\Lambda R}$, can therefore be obtained by solving Equation 1.28 with $N_{\Lambda}(d_1)$ and $N_{\Lambda}(d_2)$ calculated from separate measurements with two ^{12}C targets with different thickness ($d_1 = 1\text{ cm}$ and $d_2 = 5\text{ cm}$). Finally by analyzing $\sigma_{\Lambda R}$ in conjunction with Glauber theory discussed in Section 1.3.1, the rms radius of the hypertriton can be extracted. Simulation results [138] predict a halo rms radius for hypertritons of 4.9 fm with an interaction cross-section of $861 \pm 129\text{ mb}$, whereas the rms radius of non-halo hypertritons is estimated to be 2.8 fm with a cross-section of $645 \pm 106\text{ mb}$.

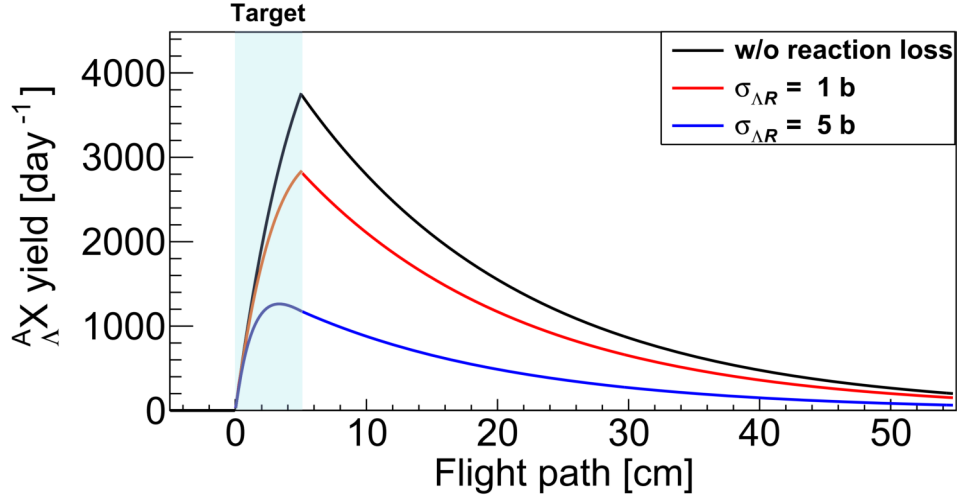


Figure 1.15.: Yield distribution of hypernuclei $^A_{\Lambda}X$ with different interaction cross-sections $\sigma_{\Lambda R}$. The ^{12}C beam intensity is 10^6 pps . The target thickness $d = 5\text{ cm}$ (blue shaded band). The figure is reprinted from [138] under CC BY 4.0.

This thesis focuses on the development of the HYDRA-TPC, a time projection chamber specifically designed for tracking pions within the HYDRA experiment. Chapter 2 presents fundamental principles governing the design, construction, and operation of TPCs. Chapter 3 provides a comprehensive overview of the experimental setup and delves into the design and construction of each individual component for the HYDRA-TPC. Calculations of the drift electric field for the TPC and simulations of electron drift are subsequently detailed in Chapter 4. Chapter 5 details the characterization of the TPC, including gain measurements. Chapter 6 presents the conduction and analysis of dedicated laser measurements with and without an applied magnetic field. Finally, the thesis ends with a summary of conclusions and future perspectives.

2. Physical Background and Basic Concepts for Time Projection Chambers

2.1. Interaction of Radiation with Gases

2.1.1. Interactions of Heavy Charged Particles with Gases

While traversing through a gas medium, charged particles continuously undergo elastic and inelastic collisions with gas atoms and molecules through electromagnetic interactions. In this section, the discussion is mostly focused on the interaction between relativistic (velocity parameter range $0.1 < \beta < 10$) low-mass charged particles, such as pions and protons, and argon-based gases.

Main Processes

In elastic scattering, both the energy and the momentum are conserved. Given the substantial mass difference between gas atoms or molecules and incoming particles, the energy transfer to the scattering center is typically negligible. However, a small deflection is added to the incident particle's trajectory after scattering. For single scattering, this deflection can be estimated with the differential cross-section $\frac{d\sigma}{d\Omega}$ at the angle θ of the Rutherford scattering model:

$$\frac{d\sigma}{d\Omega} \propto \frac{z^2 Z^2}{m^2 \beta^4} \frac{1}{\sin^4(\theta/2)}, \quad (2.1)$$

where z represents the charge of the incident particle, Z represents the atomic number of the target, m denotes the mass of the incident particle, and β denotes the velocity parameter of the particle. Higher incident particle energies lead to lower elastic scattering cross-sections and a reduced probability of deflection. When a particle undergoes multiple scattering involving a significant number of interactions, its trajectory deviates from a straight path, exhibiting a zigzag pattern. Consequently, a beam of particles with identical incident energies presents a greater angular dispersion compared to the incoming beam. The overall angular dispersion after traveling a distance x can be quantified by

$$\theta_0 = \frac{13.6 \text{ MeV}}{\beta c p} Z \sqrt{\frac{x}{X_0}}, \quad (2.2)$$

where p is the momentum of the particle and X_0 is the radiation length, a characteristic feature of the material with more details in [139].

The major energy loss mechanism for incident charged particles traversing gases is inelastic scattering, including both excitation (soft collisions) and ionization (hard collisions). Gas atoms possess discrete electronic energy levels, while molecules exhibit additional vibrational and rotational energy states associated with atomic motions. When a charged particle interacts with a gas atom or molecule, it can transfer sufficient energy to excite an electron from a lower to a higher energy state. The subsequent de-excitation of the atom or molecule often results in the emission of a low-energy photon, corresponding to the energy difference between the two states.

When the energy transferred from an incident particle exceeds the binding energy of an electron within an atom or molecule, ionization can occur, resulting in the production of an electron-ion pair. This process is termed primary ionization. The ejected electron, possessing substantial kinetic energy, can further excite other atoms or molecules through inelastic scattering. In certain cases, the energy transferred from the ejected electron exceeds the binding energy of electrons of additional atoms and molecules, leading to secondary ionization. A small fraction of these ejected electrons acquire significant energy from the primary ionization, initiating a cascade of secondary ionizations along an extended track. Such electrons are referred to as delta electrons.

Stopping Power of Charged Particles

The stopping power of a specific type of charged particle per unit distance $S(E)$ in a specific medium with the density of ρ , the atomic number Z , and the mass number A , which is defined by the mean energy loss E per unit path length x , can be given by the relativistic form of Bethe-Bloch formula [140]:

$$S(E) = -\left\langle \frac{dE}{dx} \right\rangle = B\rho z^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\ln \frac{2m_e c^2 \beta^2 \gamma^2 W_{max}}{I^2} - 2\beta^2 - \frac{\delta}{2} - \frac{C}{Z} \right], \quad (2.3)$$

where the constant B is given by

$$B = 2\pi N_A r_e^2 m_e c^2 = 0.1535 \text{ MeV} \cdot \text{cm}^2 / \text{g}.$$

In Equation 2.3, N_A is Avogadro's constant; r_e is the classical radius of the electron; m_e is the rest mass of the electron which is $0.511 \text{ MeV}/c^2$; z is the particle charge; c is the speed of light; γ is the Lorentz contraction factor defined by $\gamma = 1/\sqrt{1-\beta^2}$ and $\beta = v/c$; δ is density correction introduced due to the polarization effect which reduces the maximum energy transfer; C is the shell correction where the binding energy of electrons in the media is included when the charged particle's velocity is comparable to electrons in shells. The medium characteristic parameter, I , describes the energy required to excite or ionize an electron from an atom or molecule of the medium. This parameter is influenced by the atomic number of the constituent atoms and can be approximated in the unit of eV with [141]

$$I = (9.76 + 58.8Z^{-1.19})Z, \text{ when } Z > 12. \quad (2.4)$$

For a compound or a mixture of media, a weighted I is given by

$$I = \exp \left\{ \frac{\sum w_j (Z_j/A_j) I_j}{\sum w_j (Z_j/A_j)} \right\}, \quad (2.5)$$

where w_j , Z_j , A_j , and I_j are the fraction, atomic number, atomic mass, and mean ionization energy of the composition j , respectively. Experimentally, the mean ionization potential is often measured together with the shell correction term for simplicity. Further details can be found in [142]. W_{max} of Equation 2.3 is the maximum kinetic energy that can be transferred to one electron, given by

$$W_{max} = \frac{2m_e c^2 \beta^2 \gamma^2}{1 + 2\gamma m_e/m_0 + (m_e/m_0)^2}.$$

For a particle much heavier than electrons ($M \gg m_e$), W_{max} is given by

$$W_{max} \approx 2m_e c^2 \beta^2 \gamma^2.$$

In the case of a target mixture or a compound target, the average $S(E)$ is calculated from the weighted summation, as given by Bragg's rule [143],

$$\langle S(E) \rangle = \sum_i w_i S_i(E), \quad (2.6)$$

where w_i denotes the mass fraction of component i .

The Bethe-Bloch formula is applicable for a wide range of $\beta\gamma$ values, typically within the approximate range of 0.1 to 1000. However, the exact range may vary depending on specific materials and particle types. The mean energy loss rate per unit path length and density for muons, pions, and protons, as calculated using the Bethe-Bloch formula, Equation 2.3, is illustrated in Figure 2.1.

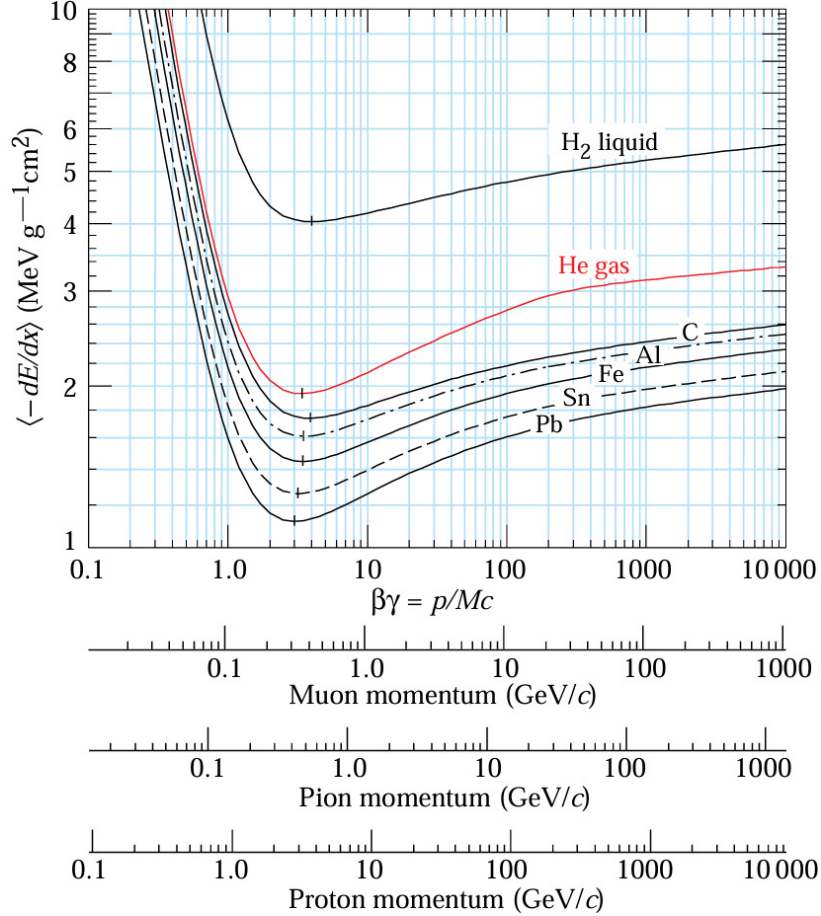


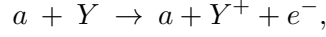
Figure 2.1.: Mean energy loss rate of muons, pions, and protons in liquid (bubble chamber) hydrogen, gaseous helium, carbon, aluminum, iron, tin, and lead. The figure is reprinted from [144] under CC BY 4.0.

As suggested both in Equation 2.3 and Figure 2.1, $S(E)$ is independent of the mass of the incoming particle, though it exhibits a direct proportionality to the square of the particle's charge-to-velocity ratio, $(z/\beta)^2$. In the low-energy regime, $S(E)$ decreases approximately like $1/\beta^2$. It reaches a minimum value where $\beta\gamma \sim 3.5$. At this point, all particles with the same charge have the same minimum energy loss rate, which only varies with the target material Z/A . Such particles are called Minimum Ionizing Particles (MIP). In the high-energy regime with larger values of $\beta\gamma$, the logarithmic term in Equation 2.3 rises, consequently elevating $S(E)$. This effect is called the relativistic rise. It is attributed to the relativistic contraction of the particle's transverse electric field, resulting in an increased interaction cross-section with the medium. However, the rise is significantly suppressed by the density correction in dense media. In the gases, the suppression is less pronounced than in solid targets. For pions at a total energy of 500 MeV, their momentum is approximately 480 MeV/c, which corresponds to the minimum energy loss in Figure 2.1. Thus, $S(E)$ of these pions in the argon gas at room temperature and the atmospheric pressure (density: 1.78 g/L) is approximately 3.1 keV/cm.

Ion-Electron Pair Production in Gases

The ionization of gas atoms and molecules results in the production of ion-electron pairs.

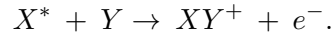
The most pronounced process is the impact ionization where a particle collides with a gas atom (or a gas molecule). The collision transfers enough energy to knock out an electron from the atom,



where a is the incoming particle and Y is a gas atom. A gas atom X may be only excited to a metastable state X^* , yet subsequently, it can interact with another atom Y and undergo the Penning ionization [145] as in the following,



or associative ionization [146], where two atoms combine into a vibrational state,



Considering all ionization mechanisms of atomic electrons, the mean free path of the incoming charged particle is given by

$$\lambda = \frac{1}{n_e \sigma_I}, \quad (2.7)$$

where n_e denotes the electron density of the gas and σ_I is the total ionization cross-section with electrons. The process is a stochastic process in which each ionization event is independent and discontinuous. The probability of production of n electron-ion pairs in a track length L follows the Poisson distribution,

$$P(k|L, \lambda) = \frac{(L/\lambda)^k}{k!} e^{-L/\lambda}, \quad (2.8)$$

where L/λ gives the number of primary ionization events. For different gases, the ionization density, i.e., the count of primary ionization events per centimeter, can be calculated as showcased in Figure 2.2.

The mean energy loss of the incoming particle for pair production in a gas is given by the W -value,

$$W = L \left\langle \frac{dE}{dx} \right\rangle / \bar{n}_I, \quad (2.9)$$

where $\bar{n}_I$ is the mean number of electron-ion pairs produced in the track length L . The W -value represents the average energy required to create an electron-ion pair in a specific gas. It accounts for all energy deposition processes, although there are other mechanisms, such as excitation, which can deposit energy in the gas. Therefore, the W -value typically exceeds the ionization potential of atomic electrons for the gas. For instance, the ionization potentials for argon (Ar) and carbon dioxide (CO₂) are 15.8 eV and 13.7 eV [147], respectively. However, the W -values for electron-ion pair production in these gases are significantly higher, at 26 eV for Ar and 33 eV for CO₂ [147].

A localized electron cluster is created when the primary electron with sufficient kinetic energy is able to ionize other electrons. The size distribution of the cluster varies, with a significant portion containing only one electron. In the case of minimum ionizing particles penetrating the argon gas, approximately 65.6% of the clusters consist of a single electron, while 15.0% contain two electrons and 6.4% contain three electrons [100]. While electron clusters can amplify signal strength, they can also negatively impact the spatial distribution of the particle track.

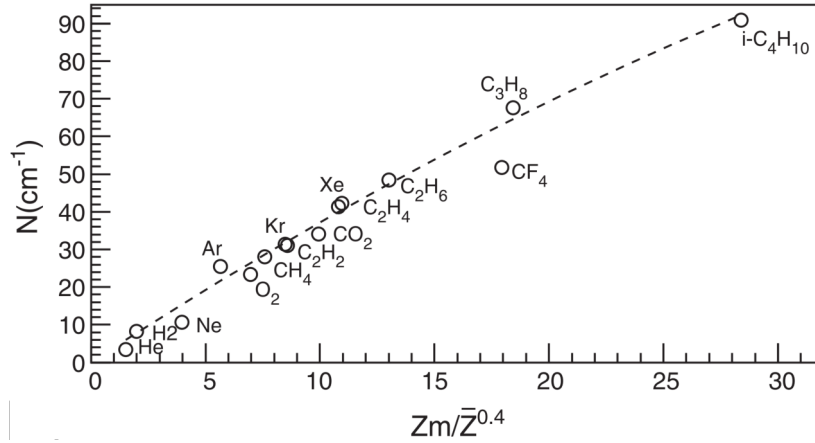


Figure 2.2.: Number of primary ionization clusters for MIPs for various gases at the room temperature and the atmospheric pressure. Z_m and $\bar{Z}$ represent the total charge and the mean atomic charge of the gas molecule. The dashed line represents fit without the point of CF_4 . Reprinted from [148], with permission from Elsevier.

The number (N) of electron-ion pairs produced in a given distance in real experiments exhibits statistical fluctuations due to the stochastic nature of particle interactions. The distribution of N is approximately Gaussian with a mean value of $\bar{n}_I$ and a standard deviation given by

$$\sigma_N = \sqrt{F\bar{n}_I}, \quad (2.10)$$

where the gas-dependent quantity F , known as the Fano factor, characterizes the dispersion of the fluctuation. For pure argon and CO_2 gases, the Fano factors have been determined to be approximately 0.2 [149] and 0.32 [150], respectively.

2.1.2. Interactions of Low-energy Photons in the Gas

Due to their lack of electric charge and rest mass, photons exhibit distinct interaction mechanisms with the gas compared to charged particles. Instead of suffering from continuous energy loss due to numerous inelastic scattering collisions within the gas, photons do not experience energy degradation. Discrete processes primarily govern photon-gas interactions, including absorption and scattering. In absorption, the incident photon's energy is completely transferred to the target atom or molecule, resulting in the photon's disappearance. For scattering events, the photon's direction is altered while its energy is either not transferred in elastic scattering or partially transferred to the target atom or molecule in inelastic scattering. As an individual photon either interacts with the gas atom or molecule in a single event or traverses the gas entirely without any interaction, the photon beam experiences an attenuation of intensity along its propagation path. The intensity (I) as a function of penetration depth (x) can be calculated with

$$I(x) = I_0 \cdot e^{-\mu x}, \quad (2.11)$$

where I_0 is the initial beam intensity; μ is the attenuation coefficient, given by

$$\mu = N \cdot \sigma_\gamma(E), \quad (2.12)$$

in which N is the gas atomic number density and $\sigma_\gamma(E)$ is the energy-dependent interaction cross-section between the photon and gas particles. For a gas mixture, the mean attenuation coefficient can be calculated

with the mass-weighted average of coefficients of constituents.

The interaction mechanism of photons with the gas varies with the photon's energy. This thesis focuses exclusively on low-energy photons in the range from ultraviolet to low-energy X-ray (below 10 keV).

Main Processes

In the low-energy regime, primary interactions between photons and gas atoms or molecules are coherent scattering and photon absorption.

Coherent scattering is an elastic scattering process where a photon interacts with a gas molecule without losing energy. Instead, the photon's direction is altered. One main type of coherent scattering is Rayleigh scattering [151]. This occurs when the wavelength of the incident photon is substantially larger than the de Broglie wavelength of the electrons in the atom. The scattering intensity is inversely proportional to λ_e^4 .

Photon absorption occurs when the incident photon's energy corresponds to the energy difference between two energy states of the gas atom or molecule, resulting in the excitation of the particle. If this energy exceeds the ionization potential, the atom or molecule undergoes ionization. The absorption spectrum exhibits a series of broadened peaks attributed to the Heisenberg uncertainty principle and Doppler broadening. Following excitation, de-excitation can occur, resulting in the emission of a photon with lower energy than the absorbed photon. Subsequent to ionization, if an inner-shell electron is ionized by photon absorption, an electron from a higher level of energy can transition to fill the vacancy. This transition releases energy, which may be in the form of photon emission or transferred to another electron, resulting in its ejection from the atom. The ejected electron is known as an Auger electron.

Ultraviolet Laser in Gases

In the ultraviolet (UV) regime of the electromagnetic spectrum, where photon wavelength ranges from approximately 100 nm to 400 nm, corresponding to photon energies from 3.1 eV to 12.4 eV. In addition to Rayleigh scattering, excitation, and single-photon ionization, a unique process called double-photon ionization can occur.

Commonly used UV laser sources like N₂ and Nd-YAG have energy of 3.7 eV and 4.7 eV. Such energies are smaller than the ionization energy of most of the gas atoms and molecules, for instance, the lowest ionization energy of 15.8 eV and 13.7 eV for Ar and CO₂, respectively [147]. Consequently, single-photon absorption does not provide enough energy to ionize the atoms and molecules.

Nevertheless, UV lasers can induce the two-photon ionization of hydrocarbon molecules. Due to the presence of numerous vibro-rotational energy states, these composite molecules exhibit a high density of vibro-rotational levels at high quantum numbers [152]. The high density allows the process for the molecule to interact with one photon and reach an intermediate virtual or real excited state that closely matches the energy of the UV laser. Subsequently, a second UV photon is absorbed before the molecule de-excites, leading to the ejection of the electron provided that the total energy of two photons exceeds the ionization energy of the molecule. For instance, toluene possesses an ionization energy of approximately 8.8 eV [153]. Although one-photon absorption using Nd-YAG lasers is insufficient to ionize toluene molecules, two-photon absorption with a total energy of 9.4 eV can effectively achieve the ionization threshold. The ionized atom density, n_i , from the double-photon ionization [152] of the ground state atom density, n_0 , is given by

$$n_i = n_0 \sigma_e \sigma_i^* N^2 \frac{T}{b} \left[1 - \frac{1}{bT} (1 - e^{-bT}) \right], \quad (2.13)$$

$$b = 1/\tau + (\sigma_e + \sigma_i^*) N, \quad (2.14)$$

where σ_e and σ_i^* are the excitation cross-section and the ionization cross-section from the excited state for the molecule; τ is the lifetime of the excited state; N is the laser flux density; T is the laser pulse width.

Experimental measurements of σ_e and σ_i^* for toluene using a 267.7 nm UV laser yielded values 6.0×10^{-18} and 2.0×10^{-17} , respectively [154]. Therefore, the ionization density, i.e., the number of electrons produced through two-photon ionization, is influenced by both the concentration of the composite molecules in the gas and the flux density of the incident UV laser. For a specific gas mixture, the ionization density exhibits a power-law dependence on the laser energy density, as illustrated in Figure 2.3.

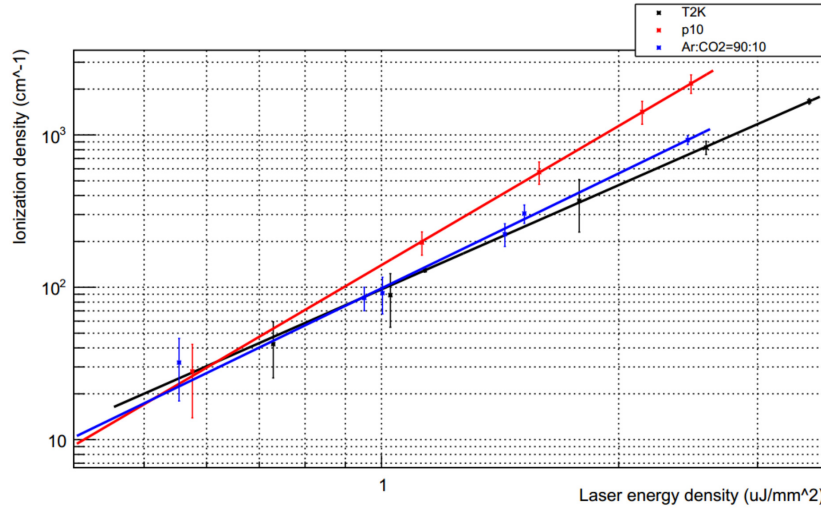


Figure 2.3.: Ionization density as a function of the laser energy density. Three gas mixtures are (1) T2K (black): Ar/CF₄/iC₄H₁₀ = 95/3/2, (2) P10 (red): Ar/CH₄ = 90/10, (3) Ar/CO₂ (blue) = 90/10. Used with permission of IOP Publishing, Ltd, from [155]; permission conveyed through Copyright Clearance Center, Inc.

2.2. Working Principles of a Gaseous Detector

Gaseous detectors are widely employed in nuclear and particle physics experiments. The construction of such detectors is relatively inexpensive due to the readily available and affordable feature of the filling gases. Additionally, their dimensions can be easily scaled to meet specific application demands without significant cost and complexity. The capability of freely adjusting the gas pressure within the detector enables a wide variety of detection scenarios. Furthermore, the low density of the gas allows charged particles to traverse a significant distance with a large acceptance angle inside the detector, offering great tracking capabilities of particle tracks with relatively high position resolution.

2.2.1. Types of Gaseous detectors

A gaseous detector typically comprises a container filled with a specific gas mixture, an anode electrode, and a cathode electrode. By applying a potential difference between these electrodes, an electric field is established within the container. The electric field separates electron-ion pairs produced in the ionization of gas atoms and molecules by charged particles, directing the electrons toward the anode and the ions toward the cathode, where they are collected for signal measurement.

There are different operation modes for gaseous detectors, depending on the voltage applied to the electrodes, as shown in Figure 2.4. The quantity of electron-ion pairs collected at the electrodes is influenced by both the primary ionization induced by the incident particle and the subsequent behavior of ions and electrons within the gas, as specified in the following operating regions:

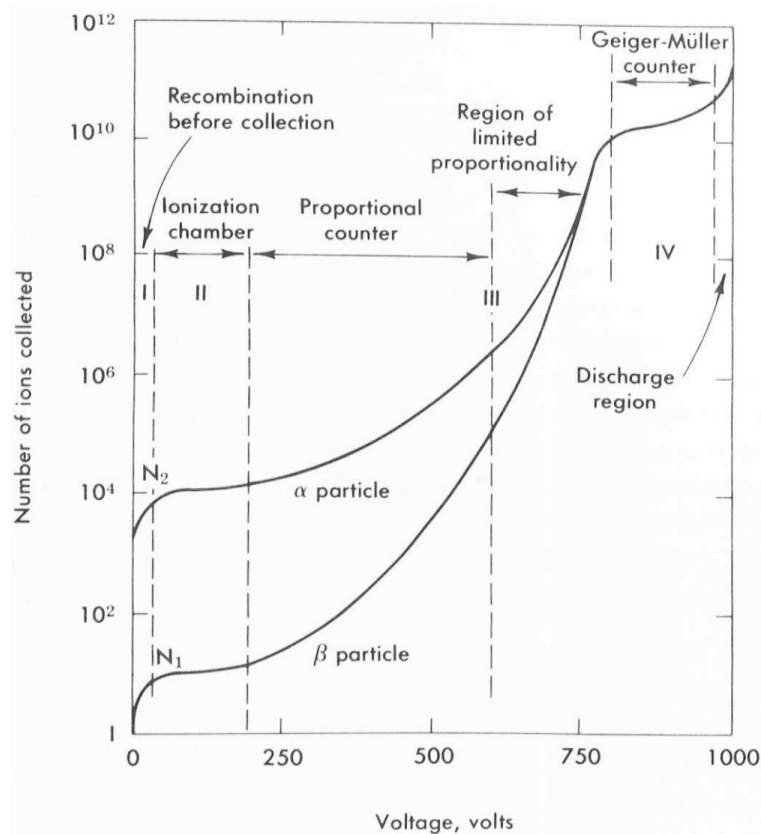


Figure 2.4.: Number of ions collected on the electrodes as a function of applied voltage in a single wire gas chamber. The figure is reprinted from [156] under CC BY 3.0.

1. Recombination region: In the absence of an electric field or the presence of a weak electric field, free electrons and ions do not experience sufficient acceleration to overcome the forces of electrostatic attraction. Thus, these electrons and ions recombine and form neutral atoms or molecules, during which energy is dissipated by the emission of low-energy photons or transferred to the newly formed atom or molecule. In this region, the signal magnitude is overly low due to the limited collection efficiency of charge carriers.
2. Ionization region: As the applied voltage increases, the recombination rate decreases and the signal rises. At a certain voltage level, almost all electrons and ions are collected by electrodes, leading to the signal saturation, corresponding to all the energy loss due to ionization. Further increases in voltage result in no effect in this region.
3. Proportional region: Once the voltage is increased beyond the ionization region, the enhanced field intensity allows primary electrons to acquire sufficient energy and initiate secondary ionization of the detector medium's atoms or molecules. This phenomenon, characterized by a cascade of secondary ionization, is termed Townsend avalanche (more details in Section 2.2.4). Within this region, the avalanche process is well controlled with a constant ratio of secondary electrons ionized per primary electron. Consequently, the magnitude of the signal exhibits a direct proportionality to the initial quantity of primary electron-ion pairs.
4. Limited proportional region: A further increase in voltage within the detector intensifies the avalanche process, leading to a rise in the magnitude of the signal. However, this amplification deviates from

the desired proportionality between the initial ionization and the resulting electrical signal. Electrons, possessing a significantly greater mobility than heavier ions, are more readily accelerated and collected by the anode leaving a cloud of positive ions accumulating in close proximity to the anode, where the avalanche process occurs. This localized concentration of the positive charge, always referred to as the space charge, distorts the electric field, which hinders the acceleration of electrons, thereby limiting the further growth of the avalanche and the linearity of multiplication. Furthermore, the distortion may lead to locally high electric field strength, which causes secondary avalanches. Further details are given in Section 2.2.4.

5. Geiger-Müller region: In this region, electrons obtain the necessary energy to not only induce secondary ionization but also significantly excite surrounding atoms or molecules, which subsequently de-excite, emitting UV photons. These UV photons are not influenced by the electric field and thereby can travel freely throughout the detector volume, where they further ionize medium atoms or molecules. Consequently, the avalanche is triggered in the entire gas volume, resulting in a saturated output signal that remains constant regardless of the incident radiation's initial energy.
6. Continuous discharge region: Sufficiently high electric fields are capable of ionizing the medium independent of external radiation sources. This leads to a continuous, self-sustained breakdown within the detector, which can result in permanent damage to the detector's functionality. Therefore, it should be avoided in practice.

2.2.2. Transport of Electrons and Ions in Gas

Diffusion in Gas

When no electric field is applied across the volume, electrons and ions created by ionization move outward in all directions due to their random thermal motion. As they spread out, they frequently collide with gas atoms and molecules, losing kinetic energy throughout the process. Eventually, a thermal equilibrium is achieved between the particles and the surrounding molecules. To simplify calculations, the probability distribution of their velocities can be approximated by a three-dimensional Maxwellian distribution [157],

$$f(v) = \left[\frac{m}{2\pi kT} \right]^{\frac{3}{2}} 4\pi v^2 e^{-\frac{mv^2}{2kT}}, \quad (2.15)$$

where m is the particle mass; v is the velocity; k is the Boltzmann's constant and T is the gas temperature. The root mean square velocity v_{rms} is thereby

$$v_{rms} = \sqrt{\frac{3kT}{m}}. \quad (2.16)$$

The diffusion equation of the particle number density n is given by

$$\frac{\partial n}{\partial t} - D\Delta n = 0, \quad (2.17)$$

where D denotes the diffusion coefficient, which depends on several factors: the particle mass, the collision cross-section between the particle and the gas molecules, the gas temperature and density [158]. Owing to their significantly greater mass, ions exhibit smaller diffusion than electrons in the same gas. The diffusion is isotropic; thus the particle distribution over the position x from the origin after time t follows the Gaussian law [159],

$$\frac{dN}{N} = \frac{1}{\sqrt{4\pi Dt}} e^{-\frac{x^2}{4Dt}} dx. \quad (2.18)$$

The two-dimensional mean squared deviation of the distribution σ_I is $\sqrt{2Dt}$. In 3-dimensional space, the value is $\sqrt{6Dt}$.

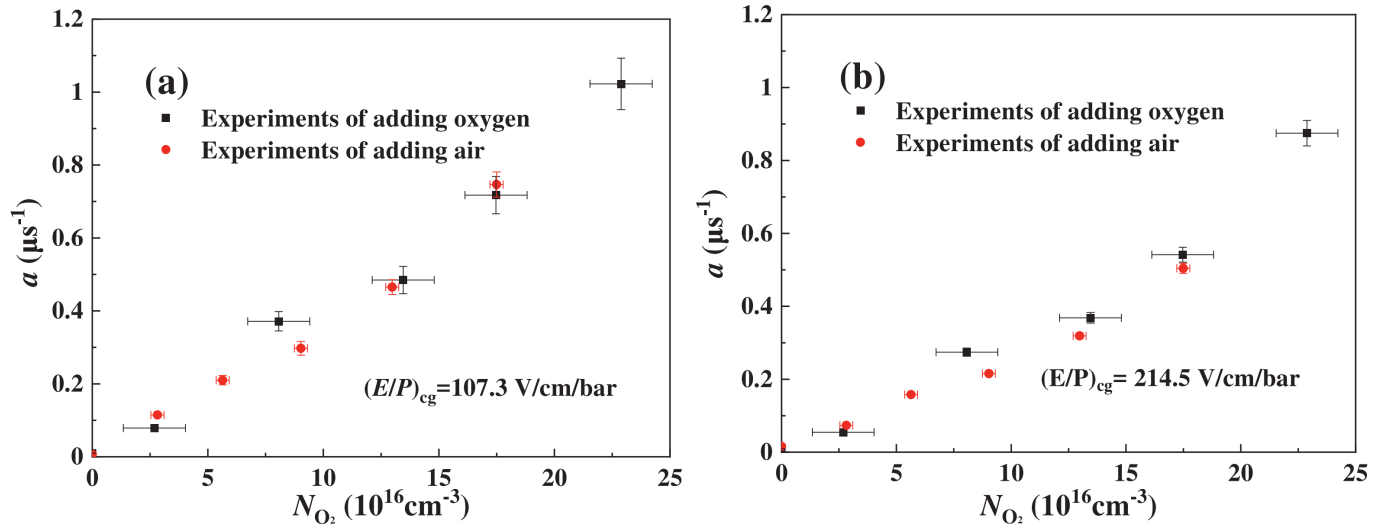


Figure 2.5.: Relationships between absorption probability a and number of oxygen molecules N_{O_2} in the operating gas, Ar + 3.45% CO₂, when the reduced electric field strength (the electric field strength over the gas pressure), $(E/P)_{cg}$, is (a) 107.3 V/cm/bar and (b) 214.5 V/cm/bar. $N_{O_2} = 10^{16} \text{ cm}^{-3}$ corresponds approximately to a fraction of 0.04% for oxygen and 0.23% for air in the gas. Reprinted from [160], with permission from Elsevier.

Drifting in an External Electric Field

In the presence of an electric field, electrons and ions experience a continuous acceleration in opposite directions. However, frequent collisions with gas atoms and molecules result in the loss of their kinetic energy. An equilibrium is established between acceleration and collisions, resulting in a net average velocity for the particles on a macroscopic scale. The overall motion is drifting towards the electrodes along the electric field lines superimposed upon the thermal random movement. The average velocity, known as the drift velocity of the charge, can be calculated with

$$u = \frac{qE\tau}{m}, \quad (2.19)$$

where q is the particle charge; m is the particle mass; E is the electric field; τ is the average time between collisions, which is inversely proportional to the collision cross-section of the particle with the gas atoms and molecules.

Thus, the drift velocity is determined by the strength of the electric field and the gas properties, including compositions, as depicted in Figure 2.6. Other factors are the gas pressure, temperature, etc. The overall influence of the gas properties is characterized by a concept called the mobility (μ) of a charged particle, defined as

$$\mu = \frac{u}{E}, \quad (2.20)$$

where u is the drift velocity and E the electric field strength. The mobility of ions is given by Einstein's relation [161],

$$\mu = \frac{qD}{kT}, \quad (2.21)$$

where q is the charge of the ion. The energy of ions obtained from acceleration is transferred by half during collisions. Therefore, the drift velocity of ions is considered to be much smaller than the thermal velocity in Equation 2.16. It depends on the reduced electric field [162], which is the ratio (E/p) of the electric field to the gas pressure and the gas temperature. Combining Equations 2.18 and 2.21, the spatial diffusion of ions is governed solely by the temperature and the applied electric field strength, while the pressure and the specific composition of the gas are irrelevant.

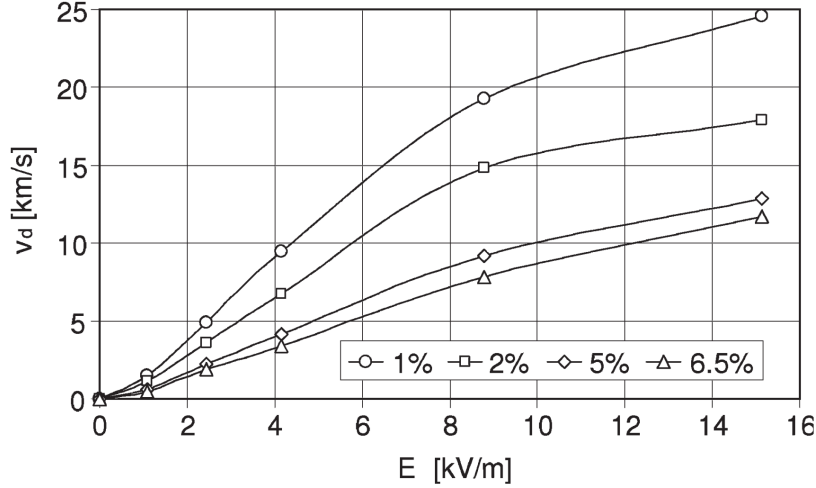


Figure 2.6.: Electron drift velocity as a function of electric field strength in Ar + CO₂ mixtures at a pressure of 1 bar and temperature. Different markers indicate the fractions of CO₂ in the various gas mixtures. © 2006 IEEE. Reprinted, with permission, from [163].

In contrast, electrons possess significantly greater thermal and drift velocities due to their much smaller mass than ions. The energy of the electron is given by

$$\epsilon = \epsilon_E + \frac{3}{2kT}, \quad (2.22)$$

where ϵ is the total post-collision kinetic energy of the electron; ϵ_E is the energy obtained from acceleration between collisions. In cold gases with many low-lying degrees of freedom, electrons experience significant energy loss in collisions with gas atoms and molecules. This energy loss can be so substantial that the post-collision kinetic energy of the electrons approaches that thermal energy ($\epsilon \approx 3/(2kT)$), leading to a similar behavior as ions in terms of their thermal motion. Therefore, the spatial diffusion of electrons in cold gases can be given by Equation 2.21 where $q = e$.

Conversely, in hot gases with fewer low-lying degrees of freedom, collisions are not effective in dissipating the kinetic energy gained by electrons. This results in a higher average post-collision kinetic energy for the electrons compared to the thermal energy of the gas atoms and molecules ($\epsilon_E \gg 3/(2kT)$). Consequently, in these hot gases, electrons can sustain a substantial drift velocity, typically on the order of centimeters per microsecond (cm/ μ s), leading to a collective motion of these electrons towards the anode. In such case, the diffusion is not isotropic anymore as the difference of diffusion coefficients given by the following [164],

$$D_L = D_0[1 - \gamma/(1 + 2\gamma)], \quad (2.23)$$

$$D_T = D_0, \quad (2.24)$$

$$\gamma = \frac{\epsilon_0}{\nu_0} \frac{\partial \nu}{\partial \epsilon} \Big|_0, \quad (2.25)$$

where D_L and D_T represent longitudinal and transverse diffusion coefficient; ϵ and ν are the mean kinetic energy of electrons and the frequency of collisions; The subscript 0 denotes the corresponding value when Δn in Equation 2.17 is zero. Figure 2.7 demonstrates the ion-like behavior of electron diffusion observed in cold gas mixtures such as 90% Ar + 10% CO₂, in contrast to the anisotropic diffusion exhibited by hot gas mixtures such as 90% Ar + 10% CH₄.

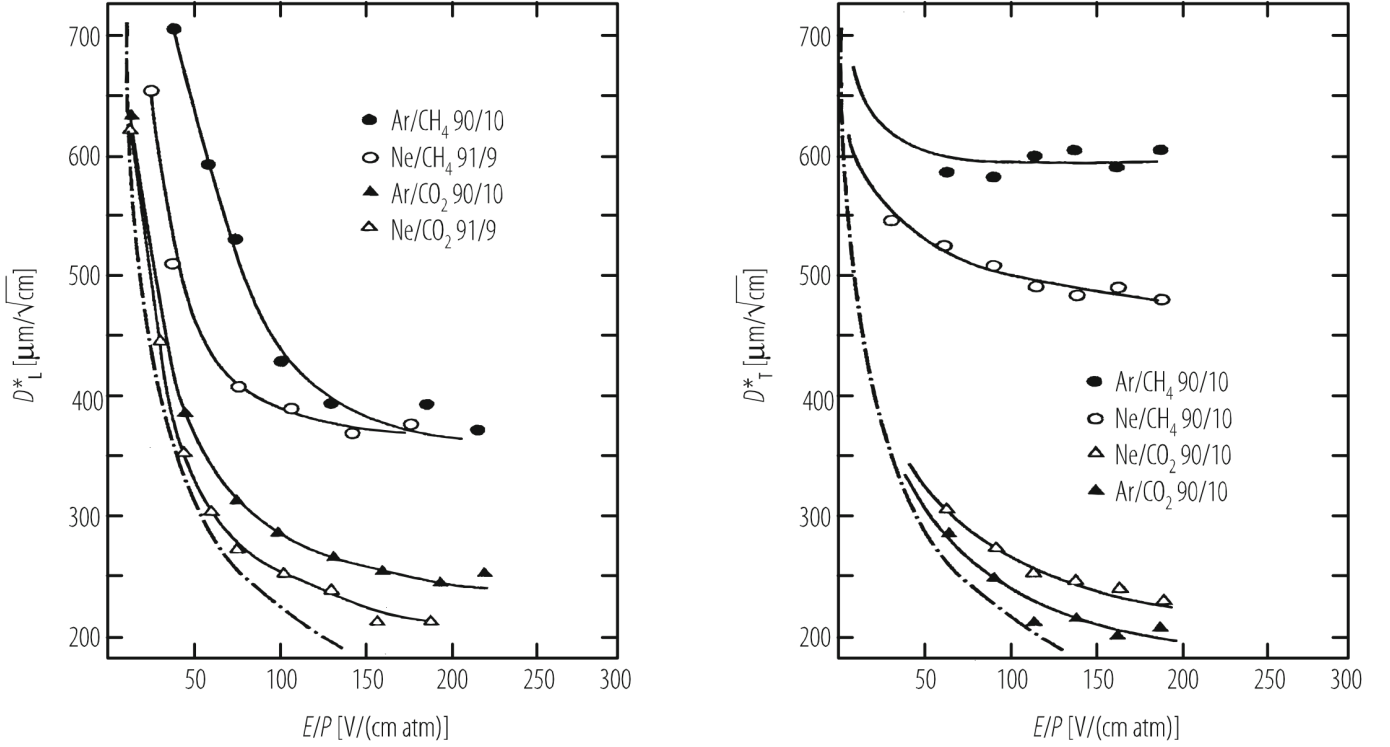


Figure 2.7.: Longitudinal and transverse diffusion constants D^* ($D^* = \sqrt{(2D/u)}$) of electrons as functions of the reduced electric field in different gases. Dash lines are calculated from Equation 2.21. The figure is reprinted from [165] under CC BY 4.0.

Drifting of Electrons in an External Magnetic Field

With the presence of an external magnetic field, the trajectories of drifting charges are altered. Within the mean free path, these particles drift in curved paths. The overall effect can be modeled as the Lorentz force superimposed under the friction force. Thus for arbitrary directions of the electric ($\vec{E}$) and magnetic field ($\vec{B}$), the motion of a charged particle (mass m and charge q) can be calculated with the Langevin equation [165],

$$m \frac{d\vec{u}}{dt} = q\vec{E} + q[\vec{u} \times \vec{B}] - k\vec{u}, \quad (2.26)$$

where $\vec{u}$ is the drift velocity of the charge; k describes the friction force strength. By defining $\tau = m/k$, the solution of Equation 2.26 is obtained as [159],

$$\vec{u} = \frac{q}{m} \frac{\tau}{1 + \omega^2 \tau^2} \left[\vec{E} + \omega \tau \frac{\vec{E} \times \vec{B}}{B} + \omega^2 \tau^2 \frac{\vec{B} (\vec{E} \cdot \vec{B})}{B^2} \right], \quad (2.27)$$

where $\omega = qgB/2m$ is the Larmor frequency with g being the g-factor ($g = 2$ for electrons). The influence of the magnetic field on the drifting of ions is negligible due to $\omega\tau \approx 10^{-10}$ [165]. For electrons, when $B = 0$,

Equation 2.26 is simplified into Equation 2.19 in which τ is the mean collision time. Two basic cases are when $\vec{E} \parallel \vec{B}$ and $\vec{E} \perp \vec{B}$. Provided that $\vec{E} = (0, 0, E_z)$ is perpendicular to $\vec{B} = (0, B_y, 0)$, Equation 2.27 is simplified to

$$u_x = -\frac{\omega\tau}{1 + (\omega\tau)^2} u_E, \quad (2.28)$$

$$u_y = 0, \quad (2.29)$$

$$u_z = \frac{1}{1 + (\omega\tau)^2} u_E. \quad (2.30)$$

where u_E denotes the drift velocity when $B = 0$. The angle between the drifting direction and the electric field, the so-called Lorentz angle θ_L is given by

$$\tan(\theta_L) = \frac{u_x}{u_z} = \frac{1}{\omega\tau}. \quad (2.31)$$

A strong magnetic field results in a small θ_L and arrival time for the drifting electrons.

In the parallel case, with $\vec{E} = (0, 0, E_z)$ and $\vec{B} = (0, 0, B_z)$, the drift velocity is the same as that in the pure electric field, given by Equation 2.19. The magnetic and electric fields may not be perfectly aligned in real-world conditions. The non-parallel alignment of the magnetic field and electric field can be evaluated through the inverse application of Equation 2.26. For instance, when $\vec{E} = (0, 0, E_z)$ is almost but not parallel to $\vec{B} = (0, B_y, B_z)$ ($B_y \ll B_z$), Equation 2.27 can be approximated as

$$u_x = -\frac{\omega\tau}{1 + (\omega\tau)^2} \frac{B_y}{B_z} u_E, \quad (2.32)$$

$$u_y = \frac{(\omega\tau)^2}{1 + (\omega\tau)^2} \frac{B_y}{B_z} u_E, \quad (2.33)$$

$$u_z = u_E. \quad (2.34)$$

In this scenario, the particle experiences a displacement after a drift length L ,

$$\Delta X = L \frac{u_x}{u_z}, \quad (2.35)$$

$$\Delta Y = L \frac{u_y}{u_z}. \quad (2.36)$$

By measuring the displacement ΔX and ΔY , the non-parallel component of the magnetic field can be determined with Equations 2.32 and 2.33.

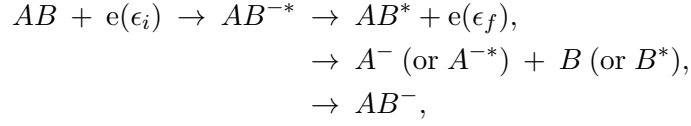
The introduction of a parallel magnetic field enhances spatial resolution. While the longitudinal diffusion is not influenced by the parallel magnetic field as suggested by Equations 2.25 and 2.34, the transverse diffusion D_T is diminished, as given by

$$D_T = D_0 / (1 + \omega^2 \tau^2). \quad (2.37)$$

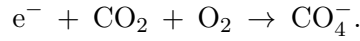
This effect on diffusion reduction is exploited to improve the spatial resolution in gaseous detectors. In a gas mixture of 90% argon and 10% CO₂ at an atmospheric pressure, employing Equations 2.19, τ is approximated as 3.98×10^{-12} seconds. Additionally, the transverse diffusion constant D_T^* in the absence of a magnetic field can be estimated to be $200 \mu\text{m}/\sqrt{\text{cm}}$ based on Figure 2.7. With the presence of a magnetic field of 2 T, the Larmor frequency of the electrons is calculated to be -3.522×10^{11} rad/s. Consequently, using Equation 2.37, the transverse diffusion constant in the presence of the magnetic field is calculated to be $67.5 \mu\text{m}/\sqrt{\text{cm}}$. For a drift distance of $L = 30$ cm, the standard deviation of the transverse diffusion is thus obtained as $D_T^* \sqrt{L} \approx 370 \mu\text{m}$.

Electron Attachment

During the transport of electrons, they are likely to be absorbed by electronegative impurities of the gas, such as the O₂ gas, creating negative ions through the following channels [166]:



where ϵ_i and ϵ_f are the initial and final kinetic energy of electrons; AB is a molecule. In these channels, the excess energy is transferred into the kinetic or internal energy of the colliding partners. The latter two mechanisms are responsible for the temporary trapping of electrons, leading to signal reduction and local distortion of the electric field. Furthermore, the presence of some quenching gas, such as CO₂ functions as a catalyst, which enhances the electron attachment effect of O₂ [167],



The number of electrons absorbed in a unit of time can be calculated with

$$N(t) = N_0 e^{-at}, \quad (2.38)$$

where N_0 and $N(t)$ represent the number of electrons from the primary ionization and at time t , respectively. The fraction of impurities in the gas directly affects the electron absorption probability, as depicted in Figure 2.5. The absorption is also influenced by both the energy of the incident electrons and the gas pressure. To minimize unwanted electron loss, the concentration of electronegative impurities within the operating gas for detectors must be kept to a minimum.

2.2.3. Amplification

In the proportional region described in Section 2.2.1, where primary electrons acquire sufficient energy from the electric field to further ionize gas atoms and molecules, electron multiplication occurs. This process is termed the Townsend avalanche, characterized by the number of new creation of electrons, denoted by dn , from n_0 primary electrons within a path length of dx ,

$$dn = n_0 \alpha dx, \quad (2.39)$$

where α is the first Townsend coefficient, which represents the collision probability per unit length. It is given by $\frac{1}{\lambda}$, where λ denotes the mean free path of the electron for a secondary ionizing collision. λ is dependent on the number density of atoms and molecules and the total cross-section for electron collision, which is dependent on the kinetic energy of electrons. The generalized form of the reduced first Townsend coefficient is given by [168]

$$\frac{\alpha}{p} = A \exp\left(\frac{-Bp}{E}\right), \quad (2.40)$$

where p is the gas pressure; E is the electric field; Parameters A and B are gas and electric field dependent. The total amount n of electrons created from the multiplication within the path length x can be calculated from Equation 2.39,

$$n = \exp \int \alpha(E) \frac{1}{dE/dx} dE. \quad (2.41)$$

In a uniform electric field, the total number of electrons is

$$n = n_0 \exp(\alpha x). \quad (2.42)$$

The multiplication factor or the gas gain G can thus be calculated by

$$G = n / n_0 = \exp(\alpha x). \quad (2.43)$$

In addition, amplification is also affected by the Penning effect described in Section 2.1.1. An enhancement can be achieved by adding specific components into a pure noble gas [169]. On the other hand, the excitation of vibrational and rotational states of molecular gases reduces the energy available for ionization, leading to quenching.

As the avalanche process progresses, the space charge effect becomes increasingly pronounced. Due to the significantly lower drift velocity of positive ions compared to electrons, they remain relatively immobile at the initial position of amplification. This relative immobility leads to a locally non-uniform distribution of charge, forming a "liquid drop" shaped region with the narrower end at the starting point of the avalanche where a high density of positive charges accumulates while electrons are more concentrated at the drop head. These positive ions can distort the electric field locally, which undermines the proportional nature of the amplification in the region.

As the electric field strength increases, the gas gain approaches the Raether limit [170], typically characterized by $G \sim 10^8$ or $\alpha x \sim 20$. In reality, the limit could be lower due to fluctuations in the avalanche [171]. When a higher electric field is applied, the concentration of electric field lines in the proximity of positive ions can initiate secondary avalanches, potentially leading to a breakdown of the gas.

2.2.4. Discharges

While amplification is essential for achieving a measurable signal, uncontrolled amplification can lead to a phenomenon known as discharge, which is a sustained flow of electrical current through the gas after breakdown. Discharge should generally be avoided in gaseous detectors, as it can limit the detector's performance by preventing proper signal collection and inducing dead time. Furthermore, repeated discharges can deteriorate the detector electrodes and insulators. To operate gaseous detectors safely, especially those working in proportional regions, one critical aspect, namely the prevention of potential discharge, should be investigated thoroughly in advance.

Townsend Breakdown Theory

As discussed in the preceding section, the Townsend avalanche is a fundamental mechanism initiating electrical breakdown. The onset criteria of the breakdown can be described by the Townsend ignition condition [172]:

$$\gamma (e^{\alpha x} - 1) \geq 1, \quad (2.44)$$

where γ is the surface ionization coefficient or the second Townsend ionization coefficient, which describes the number of secondary electrons created by a single positive charge. Furthermore, it also incorporates the electrons created from mechanisms such as ion and photon bombardment on the cathode and ionization by photon absorption. The term on the left-hand side of condition (2.44) represents the average number of electrons created by secondary avalanche from positive ions for a single primary electron. If this term is less than one, consecutive avalanches can not be sustained due to a diminishing number of electrons. When the condition is satisfied, a self-sustaining discharge occurs. The total number of secondary electrons N is given by [173]

$$N = \frac{1}{1 - \gamma(e^{\alpha x} - 1)}. \quad (2.45)$$

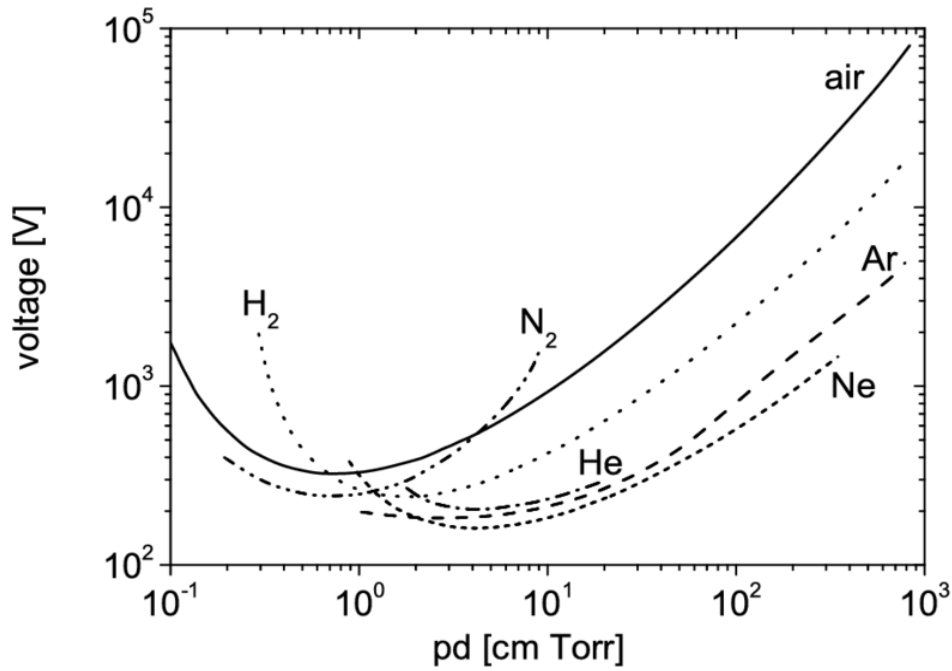


Figure 2.8.: Paschen curves, the breakdown voltage as a function of the product of pressure and distance, for different gases. The figure is reprinted from [174] under CC BY 4.0.

The breakdown voltage between two electrodes with a distance d is therefore given by

$$\alpha = \frac{\ln(1 + \frac{1}{\gamma})}{d}. \quad (2.46)$$

Considering Equation 2.40, an analytical approximation for the breakdown voltage can be obtained with

$$V_{\text{breakdown}} = \frac{Bpd}{\ln \frac{Apd}{\ln(1 + \frac{1}{\gamma})}}. \quad (2.47)$$

While A , B , and γ can be determined by measurements, at a given temperature, Equation 2.47 predicts that breakdown is determined only by the gas pressure, the potential difference, and the distance between two electrodes. The dependence of the breakdown voltage on the product of the gas pressure and the gap length is given by the Paschen curve as shown in Figure 2.8. A minimum voltage at a specific pressure exists for all gases. Below this pressure, the gas atoms and molecules are so sparse that there are insufficient secondary electrons created from ionization, thereby reducing the multiplication effect. When above this pressure, the mean free path of electrons becomes smaller; thus, electrons are more likely to lose energy before they can initialize avalanches.

Spark Discharge

In gaseous detectors, electrical discharge can occur, arising from a combination of factors, including gas composition and pressure, applied voltage, and electrode geometry. Of particular concern is the spark discharge.

After the creation of ion-electron pairs by primary ionization, secondary avalanches occur in close proximity to the positive ion tail of the avalanche liquid droplet. As consecutive avalanches propagate, they create a region of enhanced space charge towards the cathode. Finally, a streamer is formed, bridging the gaps between

electrodes and initiating a breakdown. At this stage, a decrease in the gain and detection inefficiency can be observed.

Initiated by the streamer, the spark discharge is a complete breakdown of the medium when it is sufficiently conductive between electrodes. It leads to a drastic rise in the current and thermal heating of the medium, therefore damaging the detector. The onset of the spark discharge is influenced by various factors, such as the gas mixture, electrode spacing, and electric field intensity. The gas mixture affects the production of ion-electron pairs and the breakdown voltage. The streamer length, which is the distance between the anode and cathode, also plays a crucial role, as a shorter distance results in smaller impedance and allows a higher current. For a specific detector working in the amplification region, the combined effect of the gas and electric field can be characterized as the gain. Figure 2.9 demonstrates the positive correlation between the spark probability, defined as the number of sparks per incident particle, and the gain. For the gas mixture of Ar + 11% isobutane at a gain of 4000, the discharge probability was determined to be approximately 6×10^{-6} , corresponding to an average occurrence of one discharge event per 1.67×10^5 events.

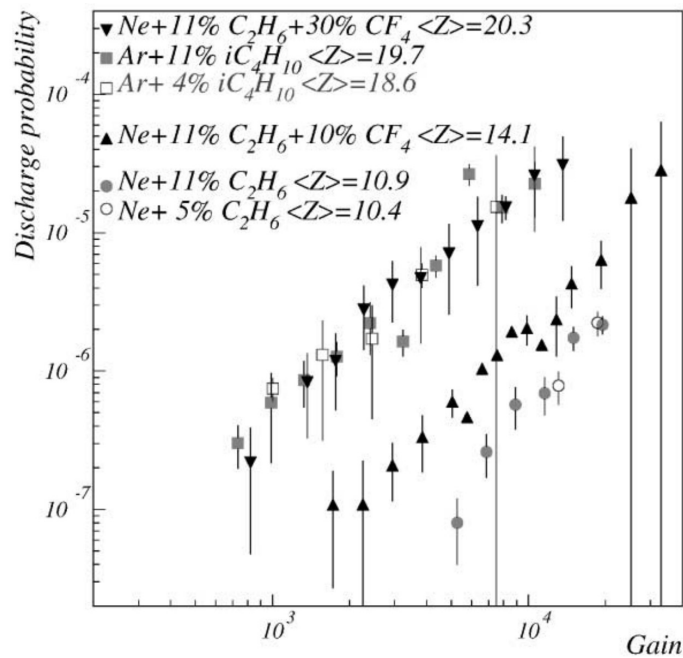


Figure 2.9.: Relation between discharge probability induced by a 15 GeV hadron beam and the detector gain on a Micromegas for various gas mixtures (more details about the Micromegas are in Section 2.3.2). $\langle Z \rangle$ denotes the mean atomic number of the gas mixtures. Reprinted from [175], with permission from Elsevier.

To minimize the risk of spark discharge, it is essential to avoid sharp edges and points on conductors. Applying surface treatments, such as coating and polishing, to conductors can help reduce field concentrations. Additionally, implementing extra shielding around the conductors can further mitigate the risk. Careful selection of the gas used in the detector is also crucial, as different gases have varying breakdown voltages. Finally, during operation, the potential difference between the electrodes should be maintained below the breakdown voltage to prevent discharge.

Triple Junction Effect

In addition to general factors of discharge, special geometry configuration can also contribute to discharges. One common configuration that exists in gaseous detectors is the triple junction. It is a joint region at the

surface of a solid insulator where the electrodes are glued and exposed to the gas. This junction divides the volume into three distinct regions: the conductive electrode, the dielectric material, and the gaseous medium. The conductor and the dielectric join at a sharp angle.

The dielectric and the gas possess different dielectric constant ϵ : ϵ of epoxy resin is typically 3.0 - 5.0, and ϵ of the gas is typically 1.0. According to Gauss's Law, electric field lines are converged in the gas region. Additionally, while the electric displacement field $\epsilon\vec{E}$ is continuous across the dielectric boundary, the electric field lines are discontinuous across the boundary, resulting in refraction and further convergence of the field lines. The discontinuity and convergence of the field lines result in a significantly higher electric field intensity than regions with smoother geometries.

Furthermore, any surface imperfections on the dielectric, such as a matte finish, can further increase the local convergence of electric field lines. Under these circumstances, a spark discharge can initiate even when the potential difference between the electrodes falls below the breakdown voltage given by Figure 2.8.

2.2.5. Operation Gas for Gaseous Detectors

In the preceding sections, the significance of working gas composition is discussed in optimizing various aspects, including ionization efficiency, electron drifting, signal amplification, and discharge stability. A typical working gas mixture for gaseous detectors often consists of the following components.

Carrier Gas

The carrier gas, which takes the main fraction of the working gas inside the chamber, ensures the uniform drift of ionized electrons. It should have low electron attachment cross-sections so that a high average drift velocity of the electrons can be achieved. A sufficiently high ionization density in the gas yields a favorable signal-to-noise ratio for accurate particle tracking while avoiding excessive space charge effects. In addition, the diffusion of drifting electrons inside the gas is supposed to be minimal.

The use of noble gases, such as argon (Ar), for most experiments is primarily driven by their full valence electron shells. This configuration results in zero electronegativity, ensuring the minimal attachment of primary ionization electrons within the gas volume. Argon possesses an ionization density of roughly 30 cm^{-1} from Figure 2.2, which is considered to be suitable for the detection of MIPs. Additionally, the non-inflammability of noble gases allows the safe operation of the detectors and their chemical inertness can minimize detector aging over time.

Quencher Gas

High electric field strengths within detectors, which are operated in the proportional region, can lead to uncontrolled amplification. The excitation of gas molecules can generate UV photons that cause the photoelectric effect on the electrodes, which initiates further avalanches and induces non-local ionization and amplification, creating spurious signals [176].

To mitigate this issue, the employment of a quencher gas is essential. A quencher absorbs photons and dissipates their energy through molecular interactions, such as vibrational and rotational transitions. Polyatomic molecules such as carbon dioxide (CO_2) possess high electron collision cross-sections for vibrational and rotational excitation within the 0.1-10 eV energy range, allowing them to effectively absorb UV photons. Figure 2.10 illustrates the electron collision cross-section by the vibration of CO_2 , which is not present for argon.

Impurity

In general, the impurity level of the working gas should be low, with a particularly low concentration of oxygen gas to minimize signal suppression caused by electron attachment. Another reason is that the oxygen gas can degrade electrode materials.

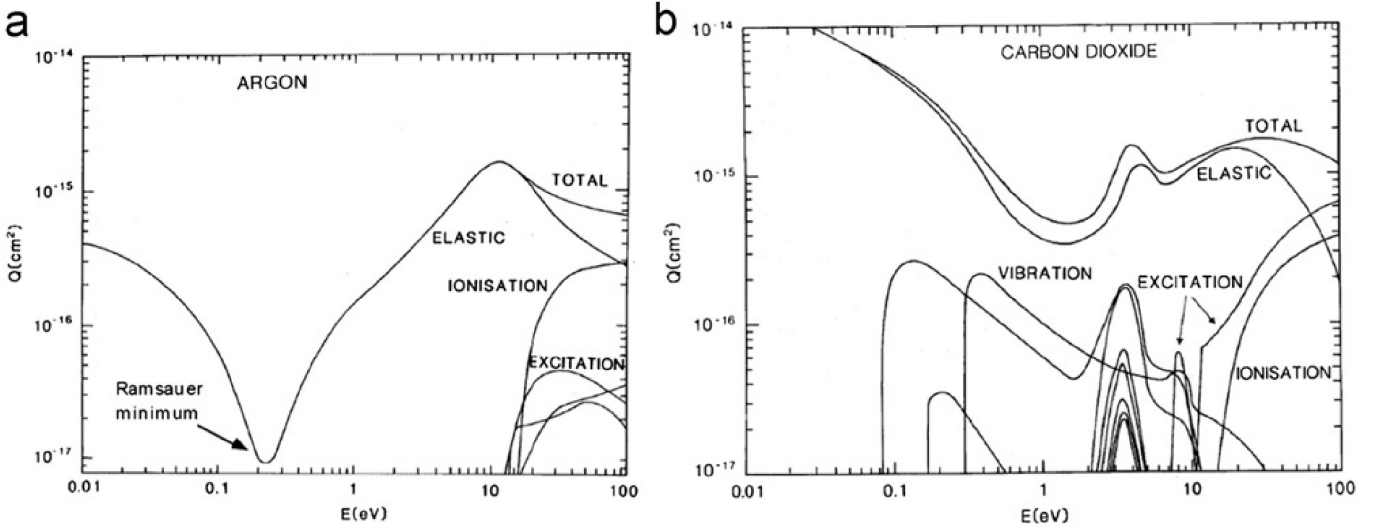


Figure 2.10.: Electron collision cross-sections in (a) argon and (b) carbon dioxide. Reprinted from [177], with permission from Elsevier.

However, some impurities can be beneficial. A small fraction of humidity, when combined with oxygen, may reduce the signal through catalytic electron absorption [167] as described in Section 2.2.2. Nevertheless, low levels of humidity (up to 100 ppmV) might improve discharge stability for gaseous detectors such as the GEM (more details are in Section 2.3.2) due to the prevention of the drying out of the glue instead of hindering amplification [178].

Furthermore, for detectors operated with UV lasers, the concentration of organic composites can contribute to the double-photon ionization as discussed in Section 2.1.2. This is impossible for pure argon or carbon dioxide gases.

2.3. Gaseous Time Projection Chambers in Brief

The Time Projection Chamber (TPC) is a gaseous detector capable of offering high-resolution, three-dimensional tracking of charged particles traversing its volume. It can be filled with either gas or liquid, offering versatility for various experimental needs. In this thesis, only relevant concepts for gas-filled TPCs are discussed.

2.3.1. Basic Concept

A gas-filled TPC operates in two distinct regions described in Section 2.2.1: the ionization region and the amplification region. It consists of a chamber filled with a working gas, as depicted in Figure 2.11. On both ends of the chamber lie conductive electrode planes, namely the cathode and the anode, with the anode plane segmented into a two-dimensional array of pads. A field cage, positioned between two electrodes, consists of conductive wires or strips. When a potential gradient is applied to these wires or strips, the field cage generates a uniform electric field dedicated to the ionization and drifting of electrons. The amplification region is positioned between the drift field and the anode. When a charged particle traverses the ionization region, it ionizes the gas atoms or molecules along its trajectory. Free electrons produced from ionization drift along the field lines towards the amplification region. Upon entering the amplification region, these electrons undergo the avalanche process. After amplification, electrons induce charges and currents on the anode pads and are ultimately collected.

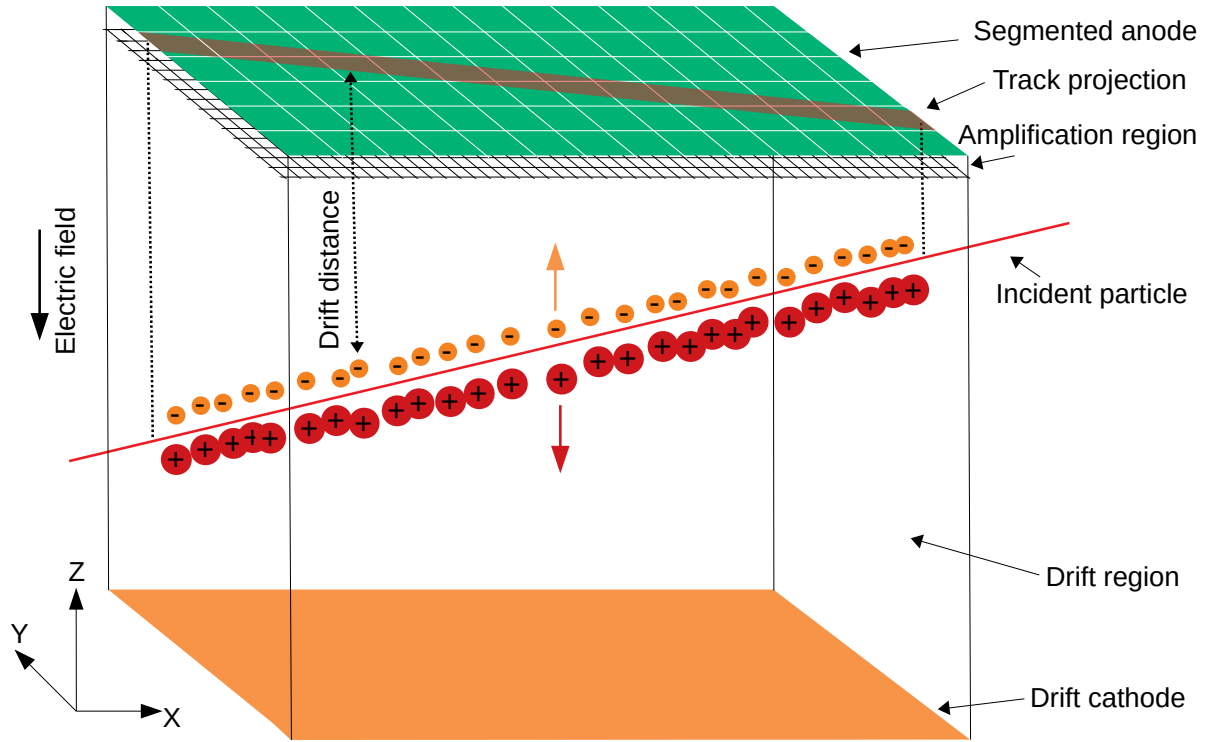


Figure 2.11.: Schematic drawing of a time projection chamber. It usually consists of the drift cathode, the drift region, the amplification region, and the segmented anode.

By analyzing the spatial distribution of the induced charge on the pad plane, the horizontal two-dimensional position of the particle track can be determined. Since the drift velocity of the electrons in the uniform electric field is relatively constant, they arrive at the anode at different times, proportional to their drift distance. Consequently, the vertical positions of the track can be obtained with the time information of each pad. Overall, the complete three-dimensional trajectory of incoming particles is reconstructed.

Originally, Multi-Wire Proportional Chambers (MWPCs) have been employed for signal amplification [179]. The MWPC is a flat gaseous detector operating in the proportional regime. A central plane of thin anode wires lies between a cathode pad plane and a mesh grid [180]. In a TPC, the MWPC is separated from the drift region by a gating grid. In the proximity of its anodes, the electric field is significantly stronger than the drift field in the field cage. Electrons undergo an avalanche when they are extracted from the drift field cage to the MWPC.

However, MWPCs suffer from inherent limitations. Mechanical and electrical constraints imposed by the wire configuration, such as wire spacing, tension, and aging, restricted the spatial resolution, which is given by $p/\sqrt{12}$ with p being the pitch between two wires. The performance of MWPCs is also limited by slow-moving ions produced from amplification. The slow mobility of ions (typically $\sim 1 \text{ cm}^2/(\text{V}\cdot\text{s})$) can lead to space charge effects that distort the electric field in the amplification region [181]. Additionally, these ions can drift back to the ionization region, causing the ion back-flow (IBF), which undermines the uniformity of the drift field. The IBF is defined as the flow of ions that are generated in the amplification region and drift back to the ionization region. The IBF fraction is dependent on the ratio of the amplification field to the drift field and the transverse diffusion of the ions [182].

The introduction of gating grids offers a solution to mitigate the issues. During the operational phase of the MWPC, when the gating grid is open, it allows electrons from the ionization region to drift freely into the

MWPC for signal amplification. Conversely, when a readout event occurs, the gating grid is then triggered and closed, effectively absorbing positive ions generated within the MWPC. However, this approach introduces dead time to the system, limiting the detector's counting rate.

To overcome the aforementioned issues, modern TPCs are mostly equipped with Micro-Pattern Gaseous Detectors (MPGD) such as Gas Electron Multipliers (GEM) [183] and Micro-Mesh Gaseous Structure (Micromegas) [184] detectors. These MPGDs offer several advantages over MWPCs, with more details in Section 2.3.2.

2.3.2. Micro Pattern Gaseous Detectors for Time Projection Chambers

Micro-pattern gaseous detectors employ microelectronic techniques, such as photo-lithography, to achieve high granularity, characterized by a pitch size that can vary from centimeters to a few hundred microns and a sub-millimeter spacing between electrodes [185]. The most applied techniques in modern physics experiments are the GEMs and the Micromegas [186].

GEM

GEMs consist of a polyimide foil, typically 50 μm thick, and a copper layer coating on both sides. These foils are etched with conical holes, with an inner diameter of approximately 50 μm and an outer diameter of approximately 70 μm , and a pitch (distance between hole centers) of 140 μm [187]. A voltage difference applied across the metal layers creates a significantly stronger electric field inside the holes compared to the drift region (as illustrated in Figure 2.12). This electric field causes electrons to undergo avalanches as they drift through the holes. One big advantage of GEMs is their flexibility. A cascade of three to four GEM foils, which significantly improve the detector stability by dividing one stage of a high gain into multiple low gain stages, or a combination of GEM foils along with other amplification detectors, such as the Micromegas, are always employed in experiments to achieve a high gain with low IBF and low discharge risk.

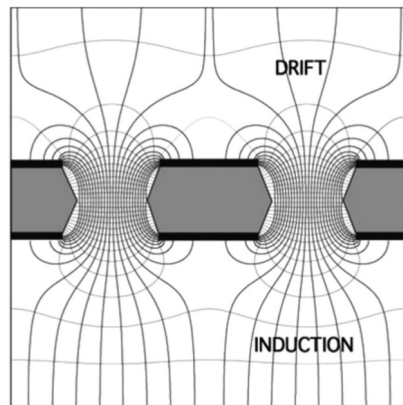


Figure 2.12.: Electric field inside holes of GEM. The induction field serves to extract electrons from the amplification towards the subsequent GEM or the anode. The figure is reprinted from [187] under CC BY 4.0.

Micromegas

A Micromegas detector is a flat parallel-plate chamber that is separated into two regions, conversion and amplification, using a micromesh that is a thin metallic grid of a few microns with an opening about tens of microns wide [184]. When electrons drift through the conversion region into the gap between the micromesh and the anode plane, the high electric field intensity (shown in Figure 2.13) inside it creates avalanches. The uniform gap of 50 to 150 μm is maintained by epoxy pillars supporting the micromesh. A more stable integration of the micromesh-anode structure can be achieved with the bulk Micromegas technology [188].

MPGDs offer several advantages compared to MWPCs. The microelectrical techniques employed in MPGD construction allow significantly smaller detector cell sizes. Consequently, MPGDs achieve spatial resolutions below 100 μm [189], while the normal resolution of MWPCs is a few hundred microns. Additionally, the time resolution of these MPGDs is typically around or less than 10 ns [190][191].

Smaller electrode structures in MPGDs result in reduced positive ion collection time and increased ion collection probabilities, leading to lower IBF rates. The IBF rate is influenced by the ratio of the amplification field to the drift field $E_{\text{amp}}/E_{\text{drift}}$ and the pitch size p . A greater $E_{\text{amp}}/E_{\text{drift}}$ ratio or a smaller p leads to a smaller IBF rate. A single-layer GEM can reduce the IBF rate to around 8% in the 90% argon and 10% carbon dioxide gas mixture with an $E_{\text{amp}}/E_{\text{drift}}$ ratio of 80 and an induction electric field of 4 kV/cm [192]. For the Micromegas with the $E_{\text{amp}}/E_{\text{drift}}$ ratio of 100, a 50 μm pitch results in a 5% IBF rate in argon-based gas mixtures, while a 17 μm pitch reduces the rate to 1% [193]. Additionally, combining MPGDs can further minimize the IBF rate. For instance, an IBF rate of less than 1% can be achieved with a hybrid detector consisting of a GEM stacked up with a 62 μm pitch Micromegas at a total gain of 1000 [194].

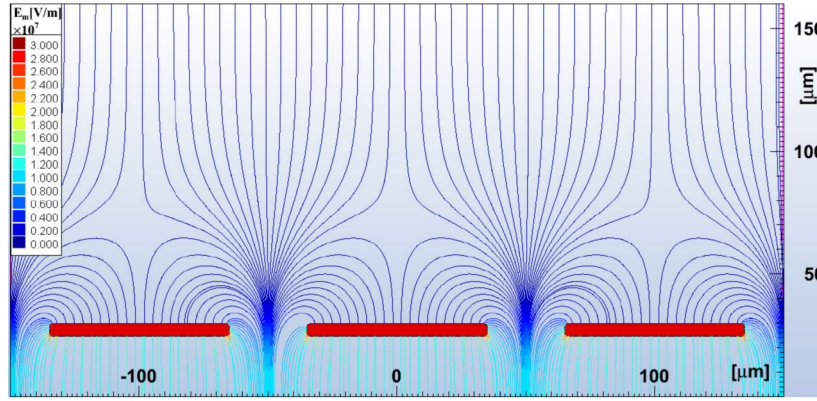


Figure 2.13.: Electric field inside holes of Micromegas with a micromesh-to-anode gap of 25 microns. The figure is reprinted from [195] under CC BY 3.0.

While MPGDs offer various advantages, they are susceptible to discharges, particularly under high voltage conditions. To reduce this risk in GEM detectors, a multi-stage amplification approach can be implemented by stacking multiple GEM foils. For the Micromegas, the introduction of a resistive layer between the readout pad plane and the micromesh serves to enhance discharge stability, especially in high-gain or high-rate applications. This layer, typically composed of Diamond-Like Carbon (DLC) with a surface resistance of a few $\text{M}\Omega$ per square, is glued onto the anode pads, offering insulating properties while enabling the application of high voltage. During streamer development within a Micromegas, the high current flowing through the DLC induces a localized voltage drop, effectively quenching the avalanche and preventing it from evolving into a destructive spark [196]. Additionally, the configuration of a resistive layer, insulated from the pad plane by a dielectric precludes the direct electrical connection between the high voltage and the electronics.

2.3.3. Signal Processing

Signal Induction on the Anode

Assuming a point charge moving towards the anode in parallel plates at an instantaneous speed $\vec{v}$, the induced current i on the electrode is given by the Shockley–Ramo theorem [197] as follows:

$$i = -q \vec{E}_w \cdot \vec{v}, \quad (2.48)$$

where q is the charge number and $\vec{E}_w$ is the weighting field at the current position of the particle. The weighting field describes the field distribution within one electrode (voltage set to 1) and any other electrodes (voltage set to 0). When an ion-electron pair is created in between the parallel plates, the induced current on the anode thereby can be calculated by

$$i = \begin{cases} e\mu^- \frac{V}{d^2} + e\mu^+ \frac{V}{d^2}, & t < t_-, \\ e\mu^+ \frac{V}{d^2}, & t_- < t < t_+, \\ 0, & t > t_+, \end{cases} \quad (2.49)$$

where V denotes the potential difference between the anode and cathode, and d their separation. μ^+ and t_+ represent the ion mobility and the arrival time of the ion at the cathode, respectively. Likewise, μ^- and t_- correspond to the electron mobility and arrival time at the cathode. Consequently, the full signal duration is given by t^+ . For instance, considering in a Micromegas, with a 128 μm gap between the micromesh and the anode pad plane, operated at a bias voltage of 400 V, the total signal collection time on the anode pad is approximately 256 ns in the gas mixture of 90% Ar and 10% CO_2 with the ion mobility of approximately 1.6 $\text{cm}^2/(\text{V}\cdot\text{s})$ [198] at the atmospheric pressure and the room temperature.

Spatial Resolution

The overall spatial resolution of a TPC for a charged particle track is influenced by a combination of factors, including diffusion in the drift and amplification region, and the anode pad size. The primary limiting factor is often the pad size, which results in an error of $p/\sqrt{12}$ where p is the pad size. Reducing the pad size can enhance spatial resolution by enabling more precise localization of the collected charge. However, the resolution decreases with excessively small pads when diffusion becomes significant. Moreover, the construction of high-granularity pads results in a substantial increase in expenses. Therefore, optimizing the detector design is crucial to achieving the desired spatial resolution.

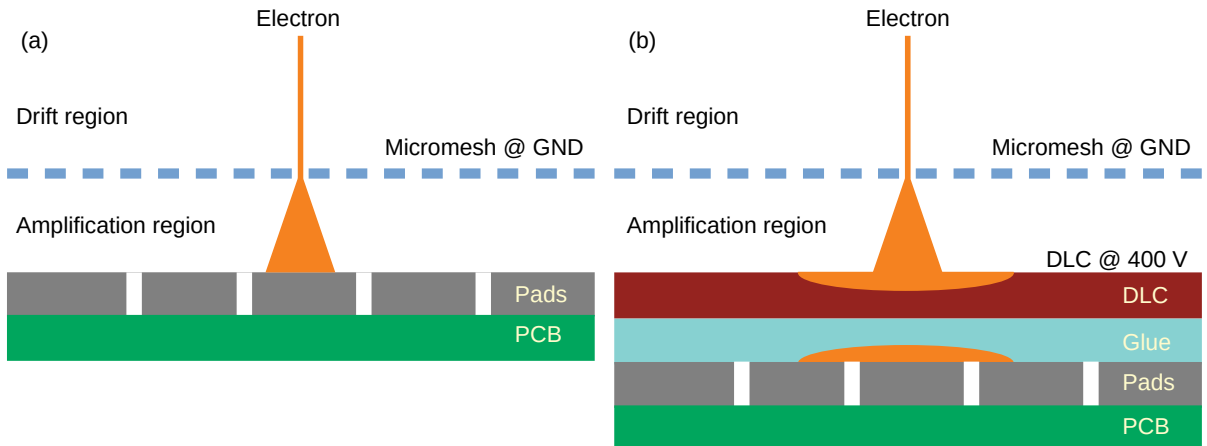


Figure 2.14.: (a) Charge amplification and collection of the Micromegas detectors and (b) charge spreading on the DLC and induction on the pads for the resistive Micromegas detectors.

Nevertheless, the introduction of a resistive layer in the Micromegas has exhibited an improved spatial resolution compared to traditional direct readout configurations [199]. The resistive Micromegas functions as an RC circuit where the resistive layer on the anode can spread the avalanche charges transversely over time, as shown in Figure 2.14. Correspondingly, charges are induced on the anode pad. The space-time evolution of

the charge density (ρ) can be calculated by the two-dimensional Telegraph equation [200]:

$$\frac{\partial \rho(r, t)}{\partial t} = h \left(\frac{\partial^2 \rho}{\partial r^2} + \frac{1}{r} \frac{\partial^2 \rho}{\partial r^2} \right), \quad (2.50)$$

where $h = \frac{1}{RC}$ is the spreading coefficient with the surface resistivity R and capacitance per unit area C . The solution to Equation 2.50 for the point charge is given by

$$\rho(r, t) = \frac{1}{2\pi h t} e^{-\frac{r^2}{4th}}. \quad (2.51)$$

This charge spreading follows a Gaussian distribution with the standard deviation of $\sigma_0 = \sqrt{\frac{2t}{RC}}$. For a Gaussian distributed charge cluster Ne with a transverse diffusion of w , the solution to Equation 2.50 is given by

$$\rho(r, t) = \frac{Ne}{2\pi(2ht + w^2)} e^{-\frac{r^2}{2(2ht + w^2)}}. \quad (2.52)$$

The corresponding deviation is $\sigma_w = \sqrt{\frac{2t}{RC} + w^2}$. For instance, considering a charge cluster drifting through a 30-cm drift region with a transverse diffusion, $\sigma_d = 370 \mu\text{m}$ in a magnetic field as discussed in Section 2.2.2, and the hybrid GEM-Micromegas amplification region with total transverse diffusion of $100 \mu\text{m}$, the final width, w , of the charge cluster arriving at the surface of the DLC layer is $\sim 383 \mu\text{m}$. The resistive Micromegas can be considered two capacitors in series, thus, the total surface capacitance is calculated by

$$C_{tot} = \frac{\epsilon_0 \cdot \epsilon_1 \cdot \epsilon_2}{\epsilon_1 d_2 + \epsilon_2 d_1}, \quad (2.53)$$

where $d_1 = 0.128 \text{ mm}$ is the distance between the micromesh and the DLC, $d_2 = 0.1 \text{ mm}$ is the thickness of the glue, $\epsilon_0 = 8.85 \times 10^{-15} \text{ F/mm}$, $\epsilon_1 = 1$ and $\epsilon_2 = 2.5$ are vacuum permittivity, relative permittivity of the gas mixture and the glue, respectively. Inserting all the values, the total capacitance per area is $5.27 \times 10^{-14} \text{ F/mm}^2$. For a surface resistance of $1 \text{ M}\Omega/\text{square}$ for the DLC and an integration time of 250 ns , the resulting deviation of charge spreading on the DLC layer is approximately 3 mm . Thus the charge is shared across multiple pads. The induced charge on individual pads can be obtained by integrating Equation 2.52 with geometrical boundary conditions [201], as

$$Q_i(t) = \frac{Ne}{4} \left[\text{erf}\left(\frac{x_{i1}}{\sqrt{2}\sigma_w}\right) - \text{erf}\left(\frac{x_{i2}}{\sqrt{2}\sigma_w}\right) \right] \left[\text{erf}\left(\frac{y_{i1}}{\sqrt{2}\sigma_w}\right) - \text{erf}\left(\frac{y_{i2}}{\sqrt{2}\sigma_w}\right) \right], \quad (2.54)$$

where $x_{i1}, x_{i2}, y_{i1}, y_{i2}$ are boundaries of the pad i . Subsequently, the implementation of charge sharing in combination with the centroid of charge calculation significantly improves spatial resolution. For a simplified 1-dimensional analysis, the centroid of charge x_c can be calculated by

$$x_c = \frac{\sum Q_i x_i}{\sum Q_i}. \quad (2.55)$$

where x_i denotes the central position of the pad i . When the aforementioned charge cluster is introduced at $\tilde{x} = 0.5 \text{ mm}$ relative to a pad plane with 2 mm pad width, where $x = 0$ corresponds to the center of a pad. The error between the cluster position $\tilde{x}$ and the reconstructed position x_c from the bulk Micromegas without a resistive layer is $500 \mu\text{m}$, while the error given by the resistive Micromegas reaches $3 \mu\text{m}$. Nevertheless, the preceding estimate does not fully describe the real-world scenario. Factors such as primary ionization fluctuations, gain fluctuations in amplification, readout electronics noise, and application of the threshold can significantly influence the overall resolution. The High-Angle TPC, equipped with a resistive Micromegas and a pad size of $9.7 \times 6.9 \text{ mm}^2$, employed in the T2K experiment, achieved a spatial resolution of $600 \mu\text{m}$ [202]. Monte Carlo simulations or direct measurements are often required to assess the spatial resolution accurately under more realistic conditions for a specific detector.

Signal Collection and Readout

Once the signal is induced on the anode, it must be efficiently amplified and shaped to ensure accurate detection and analysis. The readout electronics play a crucial role in this process, converting the induced charge into a measurable voltage signal. In modern nuclear physics, a signal readout system usually consists of the pre-amplifier, the main amplifier, the pulse shaper, the Analog-to-Digital Converter (ADC), the discriminator, the Data Acquisition (DAQ), and the terminal PC in sequence.

1. The current signal induced on the anode pads is initially weak and thereby requires amplification by a preamplifier positioned in close proximity to the detector with the intention of minimizing signal degradation due to noise and capacitance-induced interference.
2. The signal is subsequently routed to the main amplifier, where it undergoes further amplification, typically with a gain of less than $\times 1000$. The amplified signal is then processed by a pulse shaper. This component integrates the input current signal, converting it into a voltage signal. This process optimizes the signal-to-noise ratio and allows for the suppression of unwanted frequencies through the selection of an appropriate pulse-shaping time window.
3. The discriminator, positioned subsequent to the shaper, establishes a predefined threshold to effectively reject noise and signals that fall below it, thereby improving the identification of the signal pulse.
4. The last step of signal readout is the digitization of the signal with an analog-to-digital converter (ADC). The continuous analog voltage signals are converted into an n -bit binary number of 2^n bits. Subsequently, the Data Acquisition System (DAQ) performs data synchronization, formatting, and storage. Finally, the data are transferred and stored on designated devices.

The GET (General Electronics for TPCs) system [203] is an electronic platform designed to integrate all aforementioned modules. It incorporates signal amplification, shaping, and data recording capabilities. The system features a high channel density and high-speed data transfer with minimal dead time. More details are given in Section 3.3.2.

Electronic Noise

Signals are subject to degradation due to random fluctuations, commonly referred to as noise, which originate from both internal and external sources. Primary internal noise sources include Johnson-Nyquist noise, also known as thermal noise, and shot noise. Thermal noise, a consequence of random electron motion due to thermal energy, exhibits a strong temperature dependence. Its magnitude is also influenced by the circuit's impedance and amplifier gain [204]. Shot noise, arising from the discrete nature of electric current, is characterized by a Poisson distribution [205]. Its intensity is dependent on the current and amplifier gain. External noise can originate from various environmental and equipment sources. Electromagnetic interference (EMI) can be induced through inductive and capacitive coupling in detectors, electronics, and cables. Ground loops, formed when the detector is connected to separate grounds with different potentials, can introduce noise and interference due to circulating currents.

The equivalent noise charge (ENC) is a metric used to quantify the noise performance of an electronic system. It is defined as the amount of charge that produces the same amplitude of a signal in the electronics as the root-mean-square (rms) value of the output noise voltage $\langle v_{out}^2 \rangle$:

$$ENC = \frac{C_{in} \langle v_{out}^2 \rangle}{A}. \quad (2.56)$$

The rms noise is directly influenced by the amplifier gain A and the input capacitance as displayed by the ENC measurement for AFTER ASIC in Figure 2.15. To reduce ENC of electronic systems, the effective shielding of

detectors, electronics, and cables is essential. Additionally, a unified grounding system for all components is required.

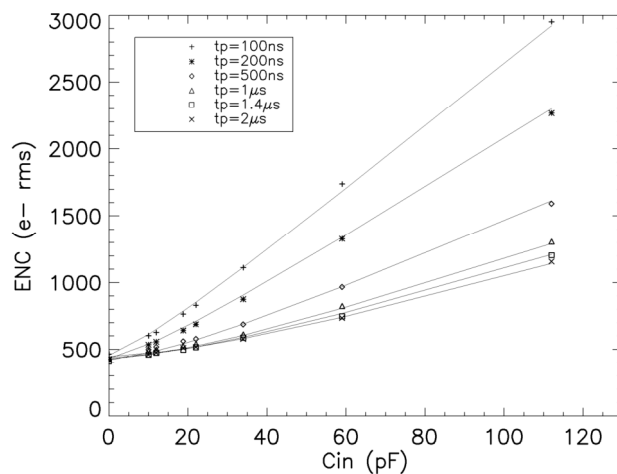


Figure 2.15.: ENC versus input capacitance for different peaking times for a fixed gain. t_p denotes the peaking time of the signal. © 2007 IEEE. Reprinted, with permission, from [206].

3. Design and Construction of HYDRA-TPC Prototype

This chapter presents the design and construction of a prototype for the full-scale HYDRA-TPC. The prototype, scaled down to one-third the size of the final detector, incorporates all essential components and technical features. Its primary purpose is to conduct a beam measurement in 2025, aimed at extracting the hypertriton radius as the first physics goal.

3.1. Experimental Setup and the TPC

3.1.1. Overview

The HYDRA experiment is proposed to be operated in the R³B terminal at GSI/FAIR (GSI Helmholtz Centre for Heavy Ion Research/Facility for Antiproton and Ion Research). The ^{12}C primary beam is produced in the ion source and sent to the linear accelerator UNILAC (UNIversal Linear Accelerator), where it gets accelerated to approximately 11.4 MeV/u [207]. Then the beam is ejected into the ring accelerator SIS18 synchrotron (Schwerionensynchrotron) where ions can be accelerated up to magnetic rigidity of 18 Tm [96]. The ^{12}C beam is accelerated to 1.9 GeV/u before it is subsequently delivered to the experimental hall of R³B (Reactions with Relativistic Radioactive Beams), which is dedicated to a large range of nuclear experiments. The whole experimental setup in the R³B hall is illustrated in Figure 3.1.

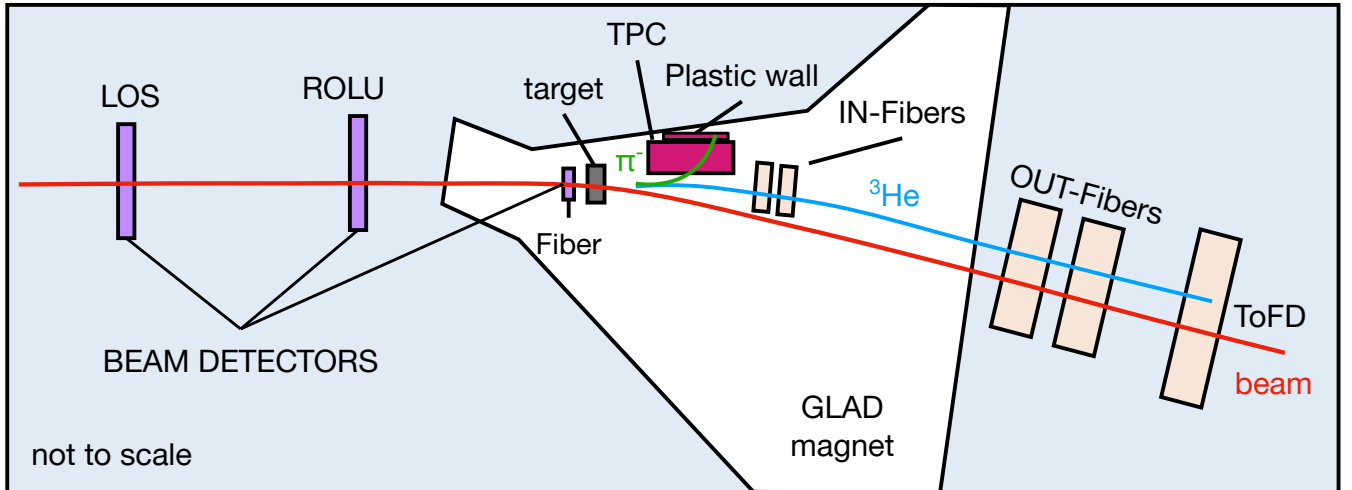


Figure 3.1.: Schematic of the experimental setup. One event of $^3_\Lambda\text{H} \rightarrow \pi^- + \text{p}$ from the $^{12}\text{C} + ^{12}\text{C}$ collision at 1.9 GeV/u is shown. The beam particle is shown in red, while the green curve is π^- and the blue curve is ^3He . The figure is taken from [208].

Beam Detectors

The incoming ^{12}C beam first traverses two beam detectors, namely LOS and ROLU. LOS is a time-of-flight detector consisting of eight plastic scintillators. It is arranged to provide the entrance beam signal, the average of all eight time measurements, used as the start of each event.

The ROLU (abbreviated from German "rechts, oben, links, unten", meaning "right, top, left bottom") detector, composed of four rectangular plastic scintillators positioned at the left, right, top, and bottom edges of a beam aperture, is employed to monitor beam centering and focus. In conjunction with LOS, it is used for beam tuning and monitoring.

GLAD Magnet and the Target

The superconducting GLAD (GSI Large-Acceptance Dipole magnet) [209] is positioned downstream of the beam detectors. The highest magnetic field flux density in the central area of the magnet is 2.2 T, corresponding to a maximum current (I_{max}) of 3583.81 A, as in Figure 3.2. As a charged particle travels through the magnet, its trajectory is deflected with a bending angle that is dependent on its rigidity. The maximum bending angle of the ^{12}C beam is simulated to be 13° .

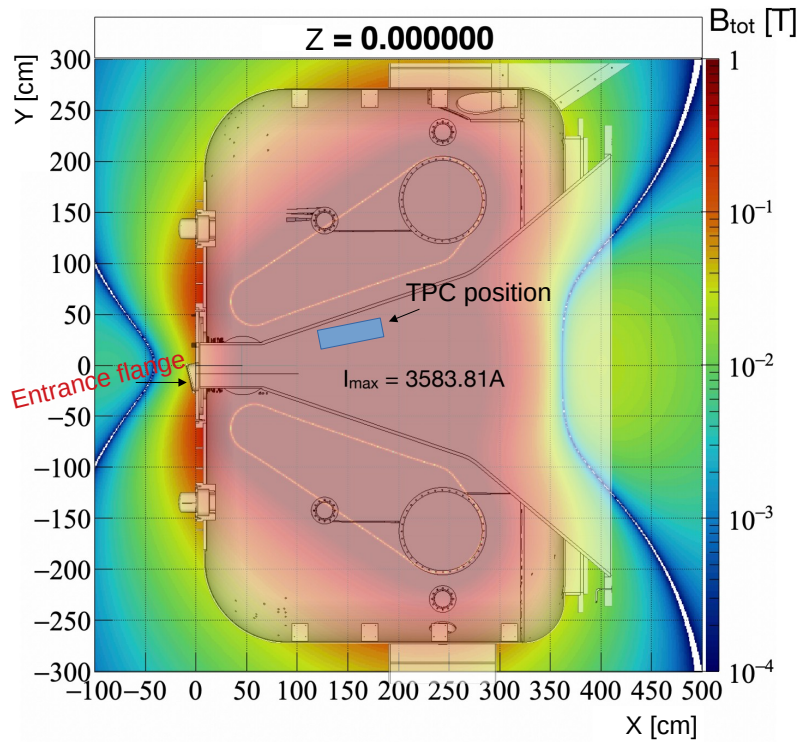


Figure 3.2.: Measurement of the magnetic field inside and in the proximity of GLAD. The color represents the ratio of the measured magnetic field to 1 Tesla, while the maximum current of the magnet corresponds to the maximum magnetic field of 2.2 Tesla. The figure is provided by Valerii Panin from GSI.

Upon entering the GLAD magnet, the beam is incident on a fiber detector and subsequently ^{12}C solid targets. As detailed in Section 1.4.2, two different targets are used in separate measurements. These targets are mounted on a remotely controlled rotary target holder, ensuring that only one target is exposed to the beam at a time. The target holder is positioned 137 cm downstream of the GLAD entrance flange. The Fiber detector ($3 \times 3 \text{ cm}^2$ wide and 1 mm thick) is placed 5 cm upstream of the target inside GLAD, which records the beam

position and ensures that the beam hits the target.

Pion Tracking Detectors

The decay products of the hypertriton's two-body decay channel, a π^- and a ^3He nucleus, diverge after the decay vertex. Due to its negative charge, the π^- meson is deflected to the left of the incident beam, as depicted in Figure 3.1, while positively charged ions, including the beam itself, are bent to the right. Consequently, these particles are detected separately. As the best compromise between the material budget and detection efficiency, the tracking of the pions is fulfilled by a combination of the HYDRA-TPC and a plastic wall at the TPC's exit window. By avoiding the direct beam, the TPC operates with a manageable space charge, enabling high-precision 3-dimensional tracking and identification of the π^- particle. The design, construction, and characterization of the HYDRA-TPC are the primary focus of this work and are detailed in subsequent chapters. The plastic wall consists of 16 scintillator bars which provide the triggering information of the pions. This detector system allows for the reconstruction of the π^- 's track and the subsequent determination of its momentum.

^3He Tracking Detectors

The positively charged ^3He traverses the magnet and a series of squared recoil Fiber detectors similar to the beam. Four IN-Fiber detectors ($13 \times 13 \text{ cm}^2$ wide and 1.5 mm thick) are positioned inside GLAD downstream of the target, reading (x,y) position signals with SiPMs, which are resistant to the magnetic field. Four OUT-Fiber detectors ($50 \times 50 \text{ cm}^2$ wide and 1 mm thick) are positioned outside GLAD, where multi-anode PMTs read the x positions of particles.

The final detector in the system is the ToFD, consisting of two frames. Each frame contains two planes of 44 vertically oriented scintillator paddles with dimensions of $1000 \times 27 \times 5 \text{ mm}^3$. Energy loss and velocity of particles are measured for charge identification. The recoil fiber detectors and ToFD are used to identify recoil particles and reconstruct their momentum.

3.1.2. Technical Requirements for HYDRA-TPC Prototype

Due to its location within GLAD, the TPC is subject to spatial constraints imposed by the magnet's inner dimensions.

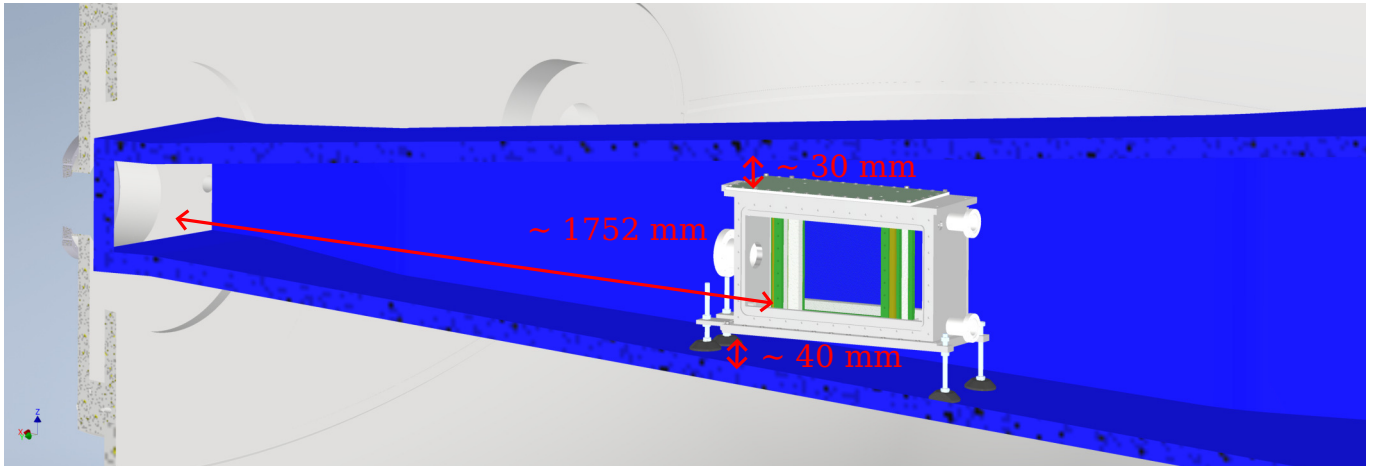


Figure 3.3.: Position of HYDRA-TPC inside the funnel of GLAD.

To optimize tracking performance by maximizing the magnetic field, the TPC should be situated in the central region of GLAD, with a leftward offset relative to the beam. However, space limitations in this area necessitate

a compromise of the TPC's dimensions thus losing the detector's detection efficiency. As depicted in Figure 3.3, the TPC is proposed to be positioned approximately 1.6 meters downstream of the entrance flange of GLAD. At this location, a π^- detection efficiency of approximately 25% [208] is anticipated.

To fulfill the momentum resolution requirement of 1% as discussed in Section 1.4.1, a spatial resolution of 200 μm in the pad plane and a drift-direction resolution of 3 mm are anticipated, as indicated by simulations [137]. Therefore, a uniform drift field is required for the TPC. Even though the approach of the charge centroid calculation method in conjunction with resistive Micromegas has been proven to improve the spatial resolution significantly, the distortion of the drift field remains a limiting factor. The distortion causes collective displacements of drift electrons that cannot be corrected by centroid calculations, thereby degrading spatial resolution. Additionally, the distortion can increase electron drift times, further contributing to resolution deterioration.

Monte Carlo simulations for a 6 cm ^{12}C target and a ^{12}C beam intensity of 10^6 particles per second (pps) predict a charged particle rate of 217 kHz inside the TPC with dimensions detailed in the following section. This rate results in a significant space charge density within the TPC's active volume due to the ion back-flow. To mitigate electric field distortions caused by the space charge, an ion back-flow of less than 1% is necessary for a gain of 4000 for the amplification region. This requirement corresponds to an average space charge density of 40 fC/cm³ in the full volume and local maxima reaching 50 fC/cm³ [210].

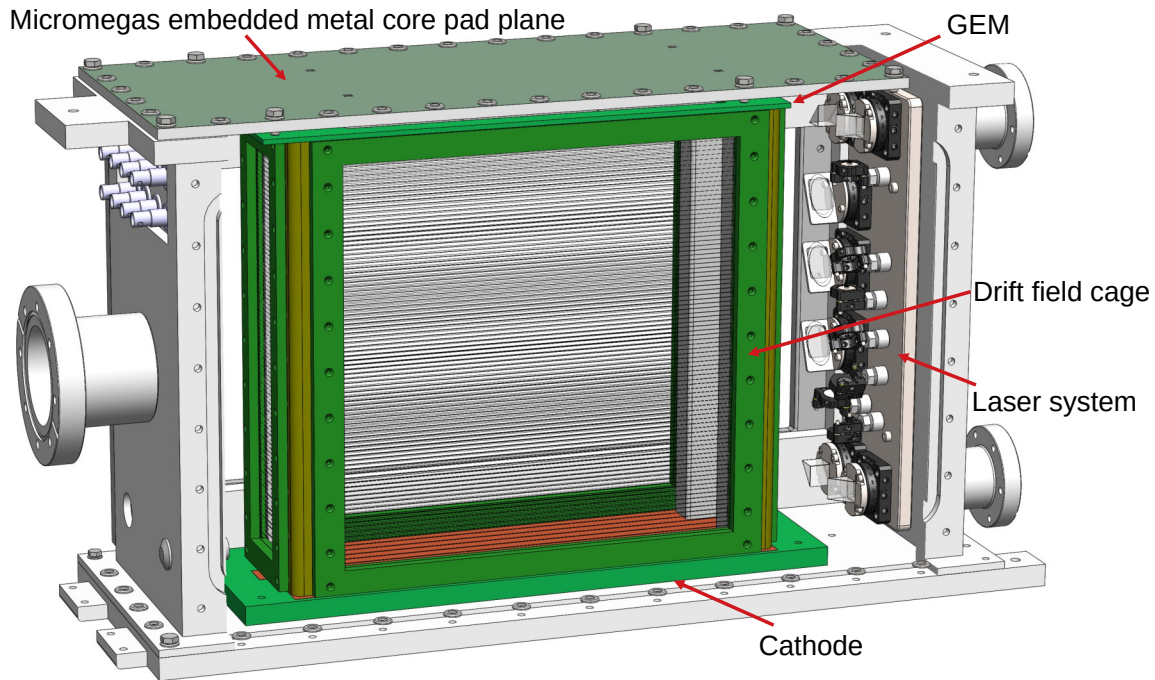


Figure 3.4.: Sectional view of the HYDRA-TPC prototype (schematic).

To accommodate these requirements, the TPC prototype is equipped with a hybrid detector of a GEM and a resistive Micromegas. A sectional view of its design is displayed in Figure 3.4. Individual components are specified in subsequent sections.

3.2. Main Components of HYDRA-TPC

3.2.1. Gas Chamber

The Vessel

A chamber, as illustrated in Figure 3.5, is built to house the drift volume, the amplification region of the TPC, and a laser system, detailed in Section 3.3.1. Its dimensions are $555 \times 229 \times 400 \text{ mm}^3$. Unless otherwise specified, a standard production tolerance of $500 \text{ }\mu\text{m}$ applies to all mechanical components in this thesis. The top plate is opened for a metal core pad plane with an embedded resistive Micromegas [211], detailed in Section 3.2.3. Two wider side walls are hollowed to create windows, allowing reaction and decay products to traverse the TPC. $12 \text{ }\mu\text{m}$ Aluminized Kapton foils are screwed onto window frames to function as windows. The chamber's front panel is equipped with eight SHV-5 connectors and a single Lemo high-voltage connector (RAD.00.113.CTM), rated for a maximum voltage of 10 kV. Additionally, this panel incorporates a Swagelok gas connector assembly comprising QC-4 Swagelok fittings of SS-6M0-61-OR, SS-6M1-PC, SS-QC4-B-6M0, and SS-QC4-D-6M0. On the rear panel, an identical Swagelok gas connector assembly is present. These two Swagelok connector assemblies are employed for flowing the gas inside the chamber. Additionally, two KF-40 flanges are welded on the back to serve as the inlet and outlet openings for the laser tracks.

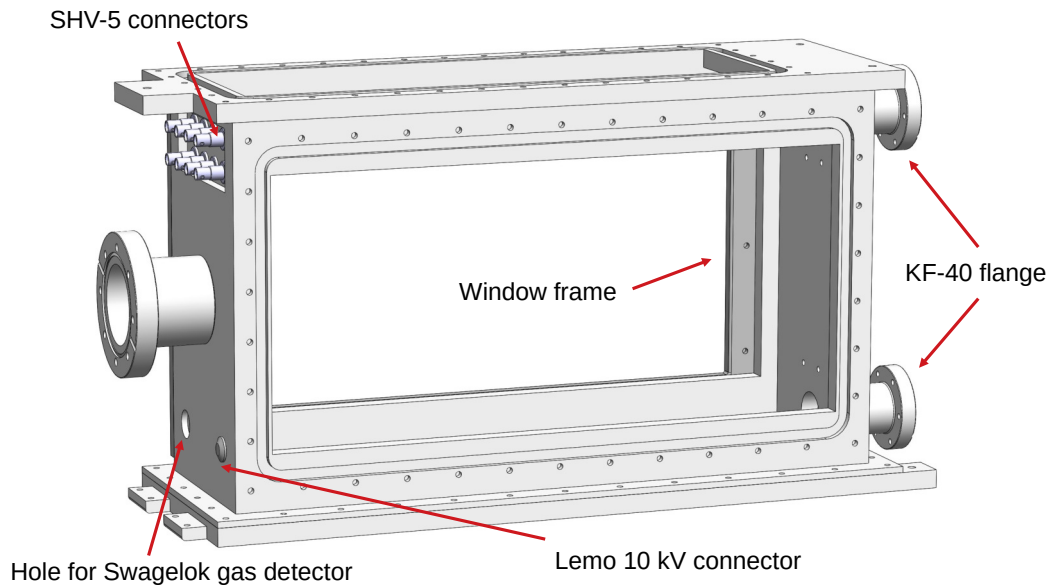


Figure 3.5.: 3-Dimensional CAD sketch of the chamber with connectors and flanges.

The chamber is constructed by welding 10-mm G.AL® C250 plates, fabricated from the non-ferromagnetic EN AW-5082 aluminum alloy, to ensure gas tightness. This material selection minimizes the effect of the external magnetic field. Ferromagnetic materials can hinder detector performance by distorting the magnetic field in the active volume. Additionally, the Lorentz force exerted on the ferromagnetic materials can lead to the detector's instability. The welding material, while initially made from aluminum, exhibits ferromagnetic properties upon heating up during the welding process. Despite this, the limited amount of ferromagnetic material introduced during welding does not significantly reduce detector performance. During measurements within the high magnetic field of GLAD, there was no significant detector instability observed.

The combined mass of all TPC components is less than 35 kg. Four feet, each with a surface area of 50 cm^2 , support the TPC, exerting a load of 0.18 kg/cm^2 on the floor of GLAD. This lightweight design allows the TPC

to be easily integrated into GLAD, which has a maximum allowable load of 0.5 kg/cm^2 .

Gas System

In accordance with the principles outlined in Section 2.2.5, the gas selection for the HYDRA-TPC is composed of argon and quencher gases. An additional requirement is that the drift velocity reaches its maximum at an electric field strength of approximately 250 V/cm , the anticipated operating field strength within the HYDRA-TPC. A potential candidate is the T2K gas mixture, consisting of 95% argon, 3% carbon tetrafluoride (CF_4), and 2% isobutane ($\text{iso-C}_4\text{H}_{10}$) with a maximum drift velocity of $79 \text{ } \mu\text{m/ns}$ at the electric field of approximately 275 V/cm as displayed in Figure 3.6(a). Near this maximum drift velocity, a broad plateau region is observed. This plateau minimizes the impact of minor fluctuations in the electric field strength on the drift velocity. The addition of CF_4 to the mixture enhances drift velocity due to its exceptionally high electron drift velocity and low electron transverse diffusion [212]. To accommodate the high charged particle rate anticipated in the final in-beam experiment, a larger proportion of isobutane or another polyatomic quencher gas may be necessary.

The selection of the optimal working gas for the final experiment requires further study, which falls outside the scope of this thesis. Instead, the 90% Ar + 10% CO_2 gas mixture is primarily utilized for the development and testing of the HYDRA-TPC due to its extensive characterization and known properties, as detailed in Chapter 2. This gas mixture exhibits lower diffusion than the T2K gas, as illustrated in Figure 3.6(b). For instance, at an electric field strength of 250 V/cm , the transverse diffusion is approximately $200 \text{ } \mu\text{m}$, whereas the T2K gas exhibits a transverse diffusion of approximately $350 \text{ } \mu\text{m}$. Additional considerations include its reasonable cost, accessibility, non-flammability, and inertness towards polymeric components within the TPC, such as certain adhesives. Nevertheless, this gas mixture is not optimal for the final experiment because of its insufficient drift velocity and the absence of a maximum value, as depicted in Figure 3.6(a).

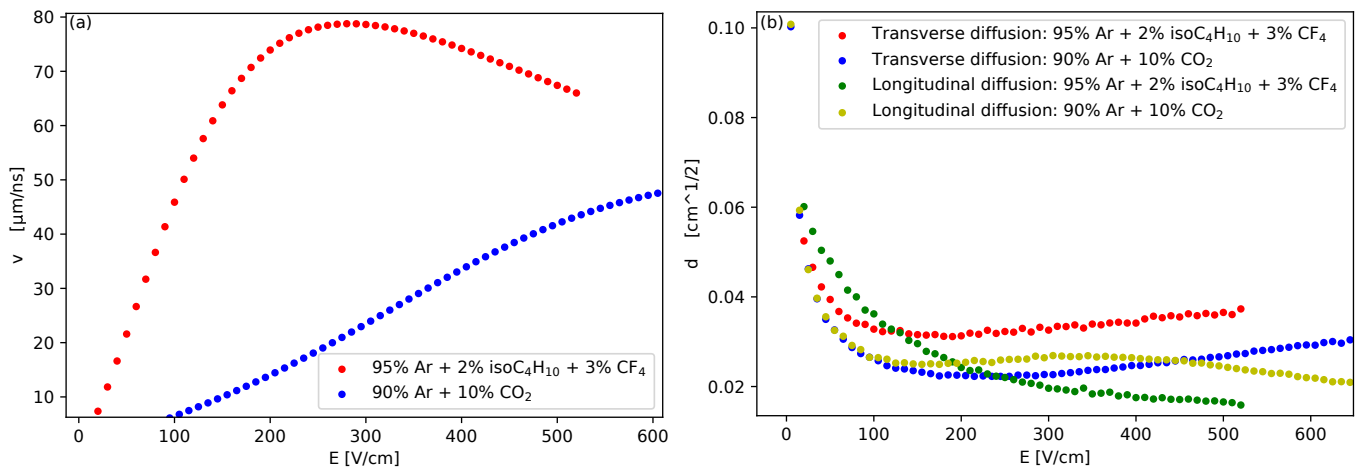


Figure 3.6.: (a) Electron drift velocity and (b) transverse and longitudinal diffusion coefficients for gas mixtures as a function of the electric field without the existence of the magnetic field. Calculation is based on Magboltz [213].

The gas chamber is connected to the gas bottle via one of the Swagelok connector assemblies, a gas pipe, and a gas regulator. The chamber's outflow is achieved through the other Swagelok connector assembly equipped with a gas pipe. During operation, the TPC maintains a steady gas flow rate ranging from 5 L/h to 20 L/h , while a slight overpressure of approximately 5 mbar is applied within the chamber. This overpressure ensures the purity and properties of the working gas by preventing the ingress of impurities through even minor leaks or permeation from the plastic components of the gas system.

3.2.2. Drift Field Cage and Cathode

A homogeneous drift field is essential for the proper functioning of the TPC. Creating such a field within a field cage while maximizing the active volume is non-trivial. Furthermore, the field cage should be as transparent as possible to minimize particle scattering during their passage. A double-layer wired field cage, inspired by the CAT [214], was designed for the prototype. To support the field cage design, simulations were conducted, as detailed in Chapter 4.

Figure 3.7 presents the field cage and its placement within the gas chamber. The field cage comprises two layers of wires carried by printed circuit board (PCB) frames, each layer consisting of two wide and two narrow PCB frames interconnected by four flexible PCBs equipped with resistors. The active area enclosed by the field cage measures $84 \times 256 \text{ mm}^2$, aligning with the pad plane area detailed in Sections 3.2.3. The vertical drift distance between the cathode and the upper edge of the PCBs is 300 mm. The minimum distance between the outermost layer of wires and the chamber's inner wall is approximately 15 mm. Based on Figure 2.8, it can be inferred that the breakdown voltage at this distance for the testing gas mixture under atmospheric conditions exceeds 10 kV, ensuring the safe operation of the experiments.

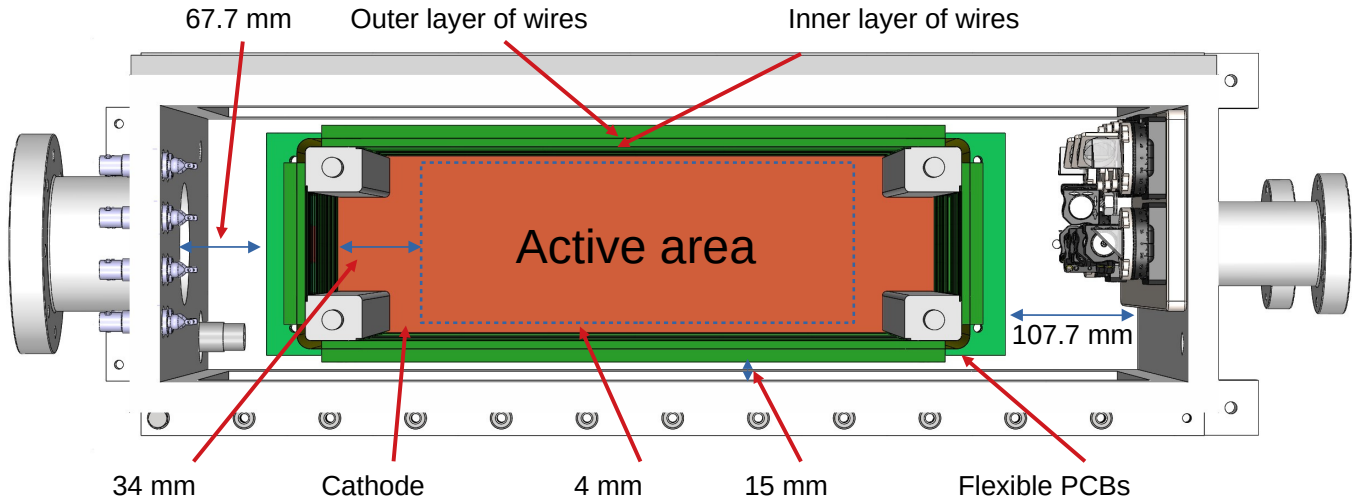


Figure 3.7.: Top view of the field cage and its position in the chamber (schematic).

The two wire layers are separated by a 4-millimeter gap, which corresponds to the thickness of the PCB frames. In each layer, the spacing between the two wires is 3 mm. These PCBs, as illustrated in Figure 3.8, serve as supports for the wires. While the PCB frames were produced by EP-DT-DD Micropattern Structures Laboratory of CERN, the wiring process was carried out using a winding machine at the GSI Detector Laboratory. The material of the wires is Beryllium Copper Alloy 25(C17200) with 98% Copper and 2% Beryllium [215], which is characterized by high conductivity, hardness, high corrosion resistance, and excellent spring properties. Its density is 8.3 g/cm^3 . The wires are tempered, age-hardened, and cold-drawn in advance, which optimizes their ductility, enabling them to hold their straight form by tension. The diameter of the wires is $75 \text{ }\mu\text{m}$. The average tension imposed on the wires is approximately 80 cN/wire.

The inner layer of wires is vertically offset by 1.5 mm from the outer layer, leading to a transparency of 94.8% for traversing particles within the window region. Apart from the wires, PCB frames also incorporate etched copper strips with the same spacing at the top, as in Figure 3.8, or the bottom side of these frames. These strips are stiffer than wires and can extend the uniformity of the drift field to both sides of the PCB frames.

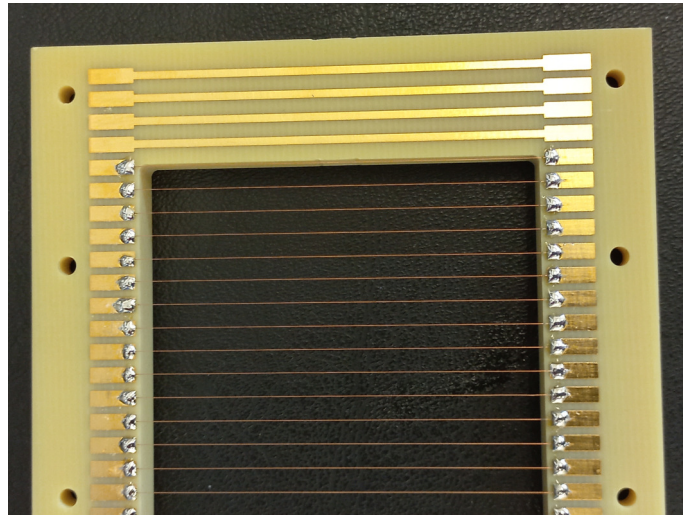


Figure 3.8.: Wiring details on the PCB frame.

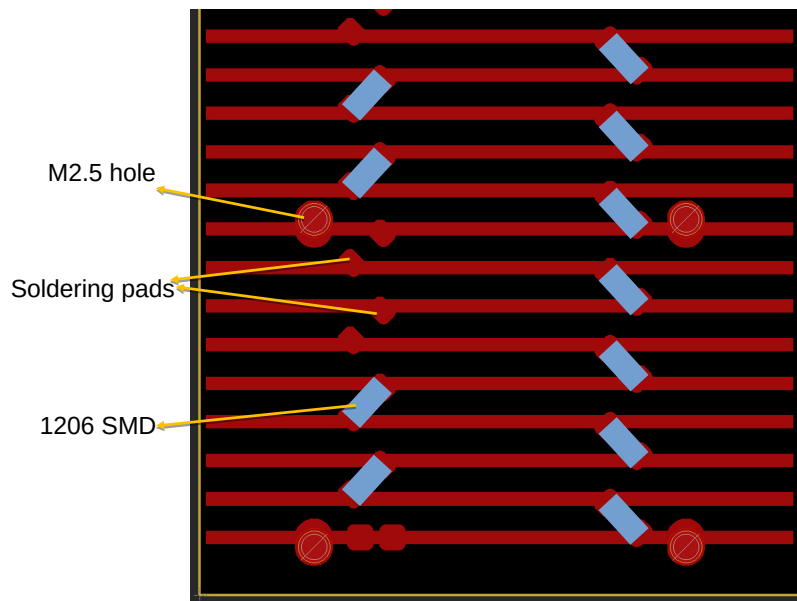


Figure 3.9.: CAD design of the flexible PCBs.

Connecting these PCB frames in each layer are flexible PCBs. These flexible PCBs, built at the laboratory of CERN, are made of 50 μm Kapton. Corresponding to the same vertical positions of the wires on the inner layer and outer layer, respectively, there are 17 μm thick passivated copper tracks with selective Ni/Au. As illustrated in Figure 3.9, soldering pads are situated on these flexible PCBs where resistors are soldered. The resistors lie between neighboring wires, forming a resistor chain that plays the central role in distributing the voltage to individual wires. Two 1 $\text{M}\Omega$ surface mounted (SMD) resistors are soldered onto the flexible PCBs between every two wires in both inner and outer layers. The power rating of resistors is 0.25 W. The first strips of inner and outer layers are connected to the cathode through 0.5 $\text{M}\Omega$ and 0.75 $\text{M}\Omega$ axial-lead resistors, respectively. The last strips are interconnected through a 0.25 $\text{M}\Omega$ axial-lead resistor and subsequently connected to an external feed-through. All the aforementioned resistors exhibit a tolerance of 1%.

To fix the flexible PCBs and wired PCB frames, four G10 pillars, with a length of 300 mm and a tolerance of 100 μm , were constructed with CNC machining. These pillars were drilled with ISO M2.5 screw holes ensuring alignment with the screw holes on flexible PCBs in Figure 3.9. Figure 3.10(a) illustrates the assembly of wired PCB frames and flexible PCBs onto the pillars using M2.5 countersunk plastic screws. Electrical contact between the wires on the flexible PCBs and the corresponding wires on the PCB frames is established through the compressive force exerted by the screws. The pillars are designed to fit into slots on the cathode PCB as depicted in Figure 3.10(b). The cathode is a 35 μm thick copper plate with an Aurum and Nickel surface finish to mitigate the triple junction effect. To prevent electrical discharge between the cathode and the chamber, which is grounded, the cathode is adhered to a 13 mm thick PCB.

The electric field strength within the drift volume is determined by the potential difference between the cathode and the top surface of the GEM (more details are in the following section). This potential difference is uniformly distributed across the resistor chain, resulting in a uniform electric field. For the test measurements conducted in this thesis, it was maintained within a range of 150 V/cm to 230 V/cm corresponding to a high drift velocity region without reaching electrical breakdown for the testing gas mixture. This translates to a potential drop of 5-8 kV at the cathode and 400-800 V at the GEM, located at the opposite end of the drift field. The total resistance between the cathode and the last strip of the field cage is 24.75 M Ω ; thus, the current estimated for a cathode voltage of 8 kV is approximately 0.3 mA.

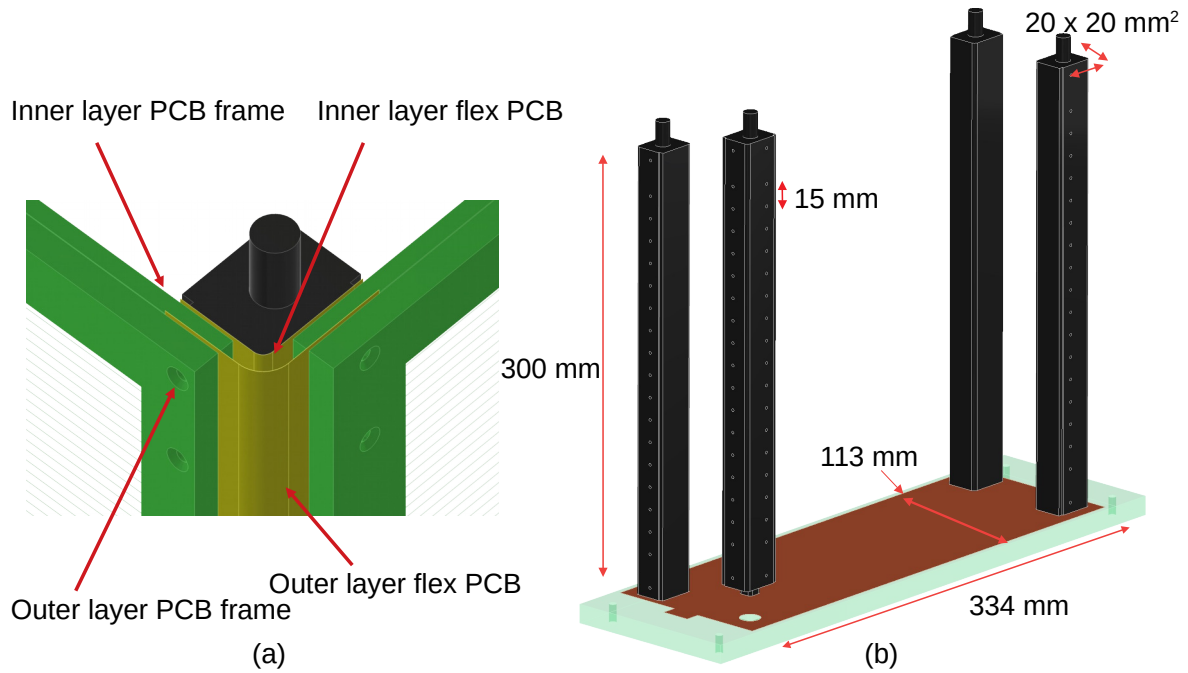


Figure 3.10.: (a) Assembly scheme of the field cage and (b) the pillars and the cathode-PCB.

3.2.3. Amplification Region and Pad Plane

Positioned above the drift field cage, the amplification region of the HYDRA-TPC is a two-stage hybrid detector incorporating a GEM and a resistive Micromegas. The Micromegas is integrated within a metal core pad plane, which houses the signal readout pads. This hybrid combination is expected to achieve an overall IBF rate below 1%, as demonstrated in Section 2.3.2.

The GEM

The initial amplification stage is a GEM foil. As depicted in Figure 3.11, it is 60 μm thick in total, with a 5 μm copper layer on each side. The GEM foil features copper and Kapton hole diameters of $70 \pm 5 \mu\text{m}$ and $50 \pm 5 \mu\text{m}$ [216], respectively. The GEM foil's bottom side is attached to a 3 mm thick PCB frame. The GEM is segmented into three segments. Each segment is electrically connected to soldering pads via resistors for both sides, as illustrated in Figure 3.11. External feed-throughs are soldered to these pads, enabling independent high-voltage application to both sides of the GEM.

The GEM PCB is positioned above the field cage, with its top surface facing the cathode. It is secured to the support pillars of the field cage using flat-head screws. The top surface of the GEM foil contributes to the uniformity of the drift field in conjunction with the cathode, the wires, and the strips of the field cage. Furthermore, the GEM foil extends beyond the edges of the active area, ensuring a uniform drift field throughout the active volume.

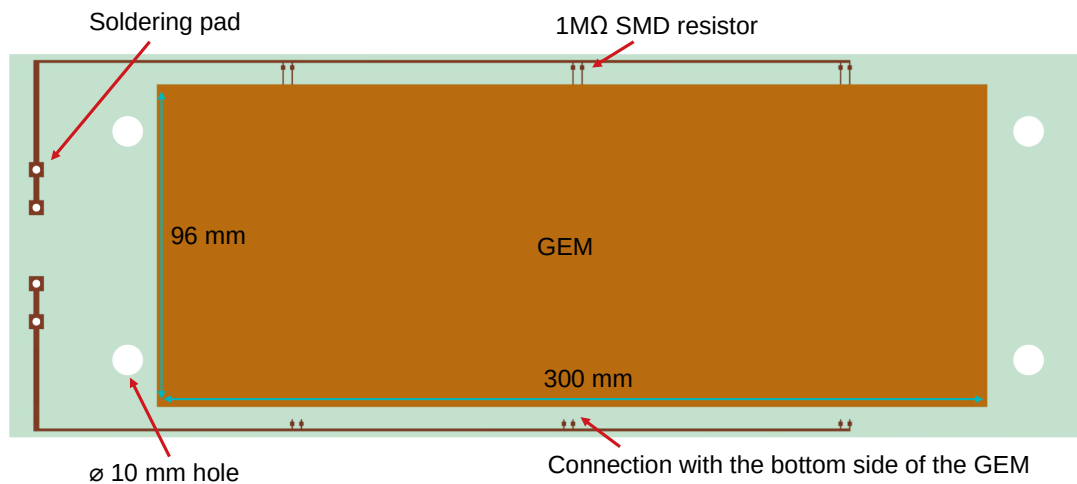


Figure 3.11.: CAD drawing for the top side of the GEM. Brown tracks denote conductive wires. The orange area denotes the GEM. The light green area denotes the PCB.

Micromegas and Metal Core Pad Plane

The second stage of the amplification region is the resistive Micromegas. A gap of a few millimeters separates the GEM and the Micromegas. Within this gap, a transfer field is applied to extract electrons from the GEM and transport them to the Micromegas. The scheme is depicted in Figure 5.1. Simulation results indicate a uniform transfer field with a 5 mm gap. Further details are provided in Section 4.2.4.

The Micromegas detector is embedded within a metal core pad plane (MCP) [211], fabricated at the CERN laboratory. This MCP structure incorporates the resistive Micromegas, SAMTEC MMTM-144-05-L-D-100 pin connectors serving as readout pads, and a bulk aluminum plate base. The primary considerations for employing the MCP include:

1. **Enhanced Mechanical Stability.** Normal pad planes, fabricated solely from PCB materials, often lack sufficient structural rigidity against mechanical deformations under significant differential gas pressure, particularly for large-area detectors. In contrast, an aluminum core within the pad plane can substantially enhance its mechanical stability, ensuring long-term stability in operation.
2. **Improved Signal Quality.** The aluminum core serves as a common ground plane, significantly reducing electromagnetic noise.

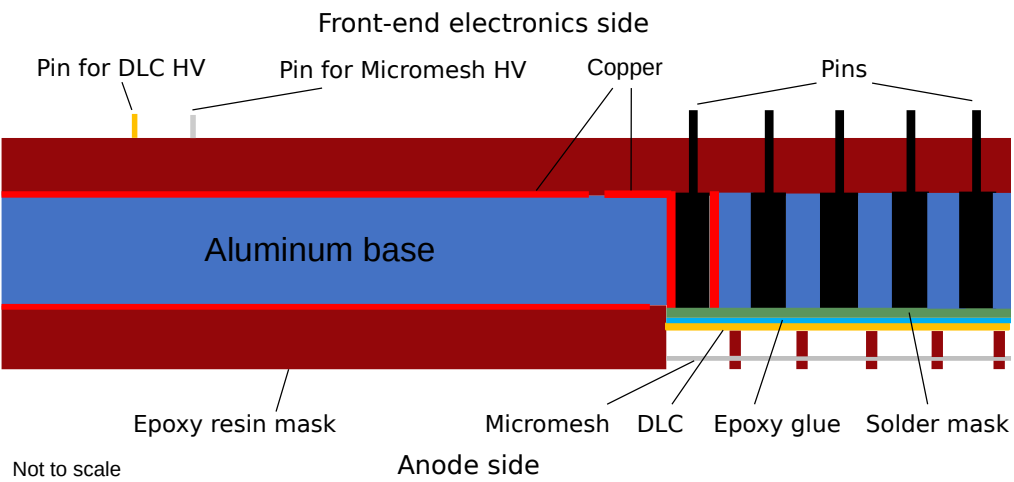


Figure 3.12.: Cross section of the metal core pad plane (schematic).

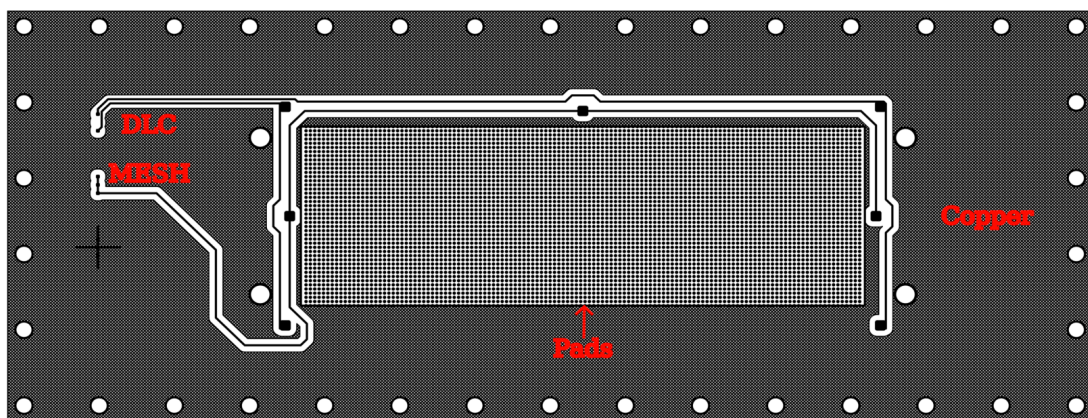


Figure 3.13.: CAD Drawing for the anode side of the metal core pad plane.

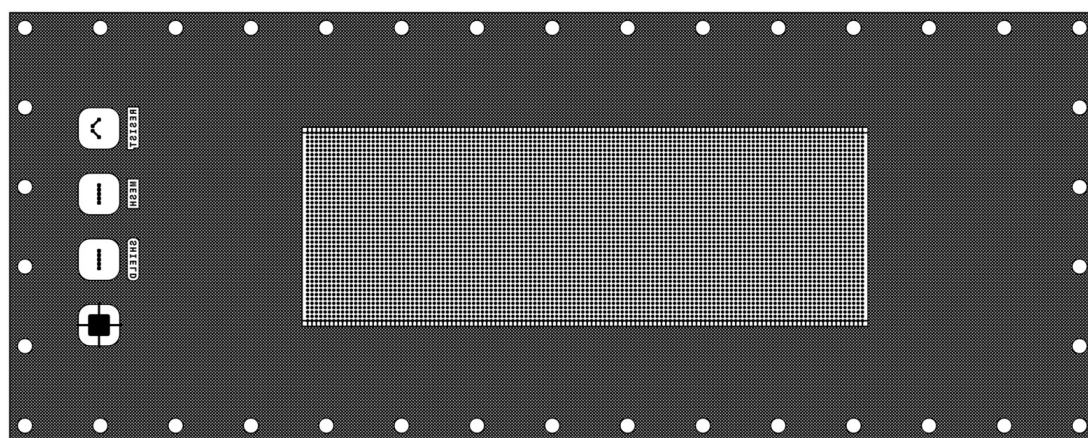


Figure 3.14.: CAD Drawing for the front-end electronics side of the metal core pad plane.

The MCPP has dimensions of $497.8 \times 197 \times 7 \text{ mm}^3$. A schematic cross-sectional view of the MCPP is depicted in Figure 3.12. The DLC layer has a resistance of $1.2 - 1.5 \text{ M}\Omega/\text{Square}$ [216]. Both the DLC and the micromesh can be connected to external feed-throughs via wires, as illustrated in Figures 3.13 and 3.14. The pad plane consists of 44×128 square pads, each covering an area of $1.9 \times 1.9 \text{ mm}^2$, separated by a gap of 0.1 mm . In total, 5632 pads occupy an area of 225.28 cm^2 , which is the active area of the HYDRA-TPC.

As Figure 3.12 demonstrates, the whole Aluminum substrate is covered by the copper area on both sides of the pad plane and can be grounded with the TPC chamber through screws. There are two rows of pads (the first row and the 44th row) connected to a dedicated copper area and subsequently shielding pins for independent grounding. They can also be directly grounded to the chamber with screws.

3.2.4. Assembly of the TPC

The following steps outline the assembly process for the main components of the HYDRA-TPC:

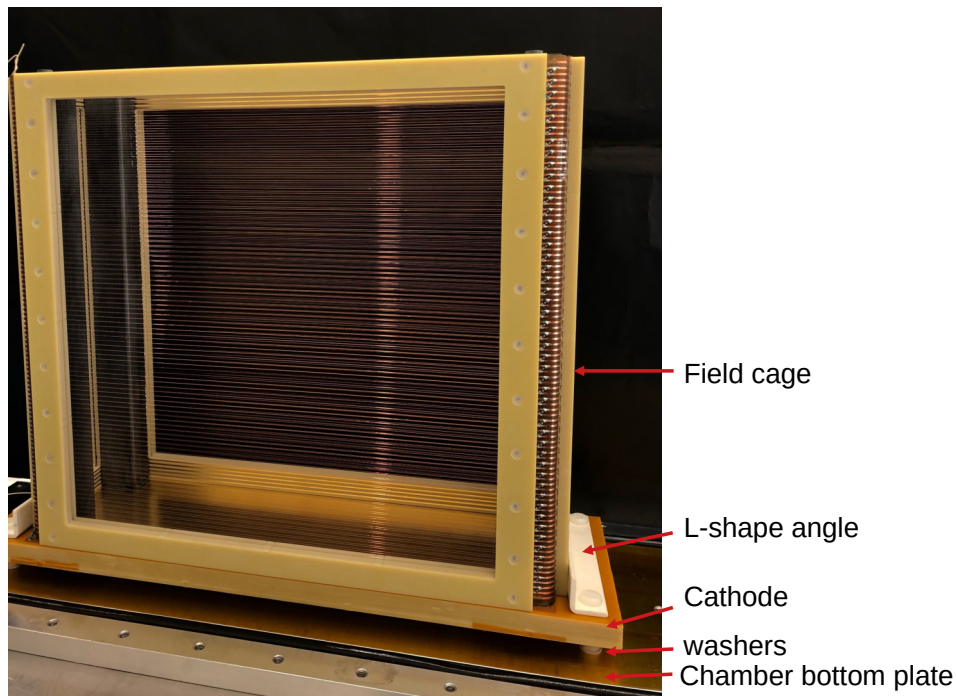


Figure 3.15.: Assembly of the field cage and the cathode on top of the bottom plate of the chamber.

1. Insert four pillars into designated slots on the cathode PCB as in Figure 3.10(b).
2. Fix the flexible PCBs and PCB frames to the pillars using Teflon screws, following the configuration outlined in Figure 3.10(a). The inner layer is fixed first, followed by the outer layer.
3. Connect the first strips of the flexible PCBs to the cathode with Allectra PEEK insulated cables.
4. Fix the PCB frames to the cathode PCB with two L-shaped plastic angles. Subsequently, the cathode assembly, separated from the bottom plate by approximately 11 mm using spacers, is screwed onto the bottom plate of the gas chamber. The assembly of the field cage, the cathode, and the chamber bottom plate is displayed in Figure 3.15.
5. Connect the last strips of the field cage to the external feed-throughs.

6. Mount the GEM onto the pillars with electrodes of both the top side and the bottom side soldered to the HV feed-throughs via PEEK cables.
7. Close the chamber with the metal core pad plane as showcased in Figure 3.16.
8. Install all gas connectors and HV connectors.
9. Install two windows on the chamber.

To mitigate dust contamination within the HYDRA-TPC, all assembly steps were conducted in the controlled cleanroom of the Detector Lab at GSI. Before being transported out of the cleanroom, the chamber was securely sealed with screws and O-rings. All spare parts were inspected with a 0.5 Tesla magnet to eliminate ferromagnetic components. Subsequently, upon removal of the TPC from the cleanroom, gas pipes were connected to the inlet gas mixture, 90% Ar + 10% CO₂. This mixture was continuously flushed at a nominal flow rate of 12 L/h to minimize the oxygen concentration within the TPC. An oxygen sensor was installed at the gas outlet of the chamber following assembly to monitor the oxygen concentration in the exiting gas flow. After a two-day flushing period, the oxygen concentration within the TPC was consistently maintained below 30 parts per million (ppm).

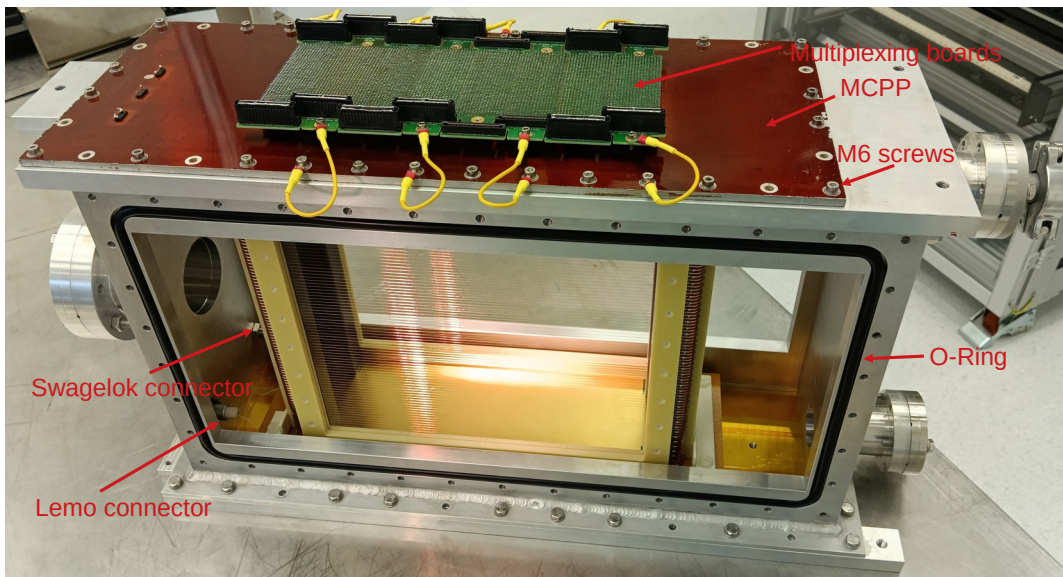


Figure 3.16.: Assembly of the HYDRA-TPC without windows. On top of the TPC are multiplexing boards, grounded with the chamber through yellow cables.

3.3. Peripheral Equipment of HYDRA-TPC

3.3.1. Laser System

The magnetic field exhibits a non-uniformity inside the funnel of GLAD. From the latest field map, depicted in Figure 3.17, the vertical component of the flux density, which is paralleled to the drift field of the field cage, has a less than 5% deviation inside the active volume. The ratios of non-parallel flux density components to the vertical component are generally less than 3%.

To evaluate the impact of the non-uniform magnetic field, a laser system was implemented. This system comprises a laser source and a mirror system with reflective mirrors and three bundles of micro-mirrors. The laser source, a 266 nm VironTM diode-pumped Q-switched Nd:YAG laser, with a nominal power output of 9.37 mJ. The laser beam exhibited a Gaussian distribution with a mean intensity of 2.48 mJ/mm².

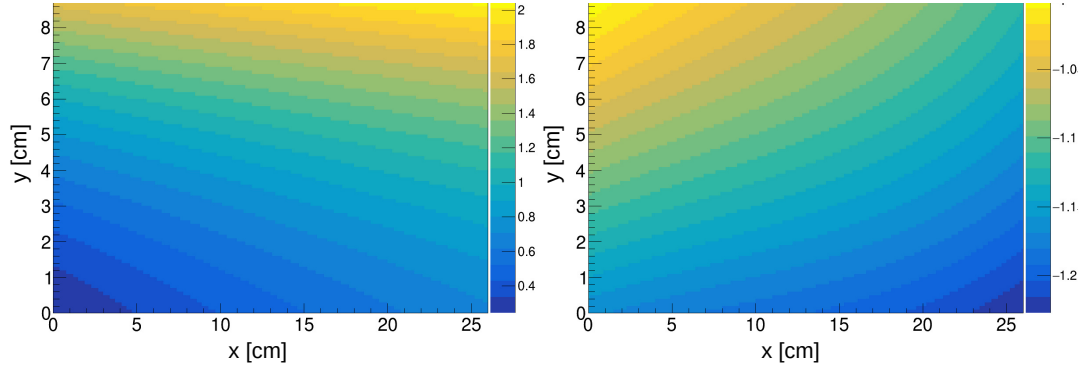


Figure 3.17.: (a) Ratio of y component to the z component and (b) ratio of x component to the z component of the flux density in the active region when The TPC is positioned inside GLAD as Figure 3.3. The percentage is represented by variations in color. The xy plane is parallel to the cathode plane, while $z = 23$ cm is the distance from the cathode. The field map of GLAD was calculated by José Luis Rodríguez-Sánchez from Universidad de la Coruña.

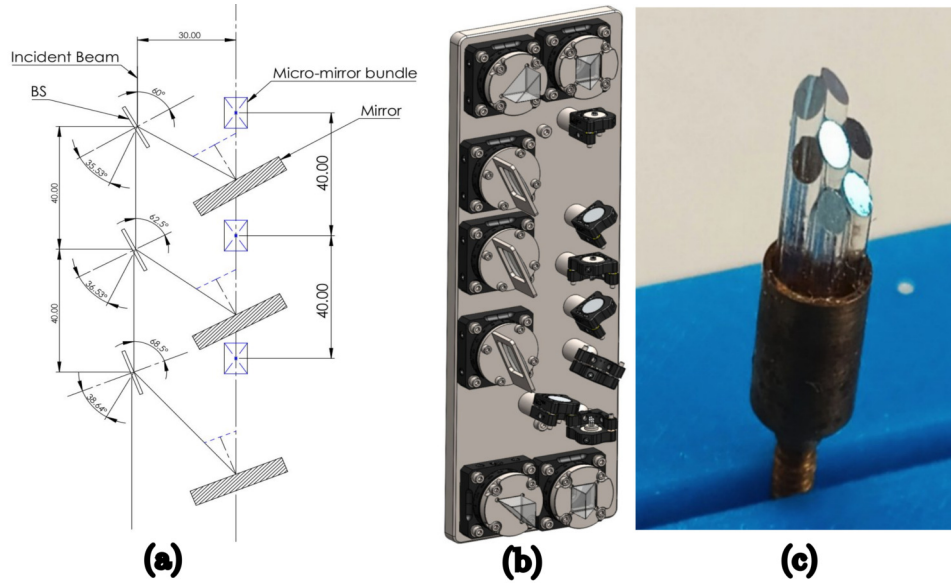


Figure 3.18.: (a) Schematic and (b) 3D CAD drawings of the mirror system. The system was built by Alexandru Enciu at TU Darmstadt. (c) A micro-mirror bundle, borrowed from the STAR collaboration. Each micro-mirror has a diameter of 1 mm perpendicular to the laser beam direction. The entire bundle has a diameter of 3 mm. Along the laser beam direction, the total length of the micro-mirrors is 5 mm. New micro-mirror bundles for the HYDRA-TPC are being constructed.

The laser source operates at a maximum pulse repetition rate (PRR) of 20 Hz, with a pulse width of 7 ns. Upon activation, the laser beam is directed through a quartz window mounted on the KF-40 flange into

the HYDRA-TPC. A series of reflective mirrors and fused silica beam splitters (BS) of the mirror system, as depicted in Figures 3.18(a) and (b), subsequently splits the laser beam into three sub-beams, each 17% of the primary beam. These sub-beams are then projected onto micro-mirror bundles, each consisting of seven 1-mm diameter micro-mirrors oriented at different angles, as illustrated in Figure 3.18(c). The micro-mirrors reflect the laser beams in various directions, generating laser tracks within the active volume of the field cage. These laser tracks produce primary ion-electron clusters through the two-photon ionization of gas molecules, which enables the calibration of drift velocity and the characterization of the magnetic field's impact on electron drift. By measuring the time difference between signals from distinct tracks, the drift time of electrons can be inferred. Combining the relative positions of the micro-mirrors within the entire drift volume, the drift velocity can be calculated. Subsequently, the non-uniformity of the magnetic field can be assessed by comparing the displacement of tracks in the presence and absence of the magnetic field.

3.3.2. Signal Readout System

The high pad plane density necessitates the use of high-density readout electronics. Within the duration of this thesis, a combination of the GET-based electronic system and multiplexing boards was utilized for test measurements. However, this configuration exhibited limitations regarding the maximum event readout rate. To address this constraint and accommodate the anticipated trigger rate of approximately 30 kHz in the upcoming beam experiment [210], the high-rate VMM3a front-end electronics integrated with the Scalable Readout System (SRS) [217] is planned for deployment. Nonetheless, this thesis focuses solely on measurements conducted with the GET system.

GET-Based Electronic System

The GET (General Electronics for TPCs) system [203] is a high-integration electronics system that incorporates amplification, signal processing, and data acquisition in a single unit.

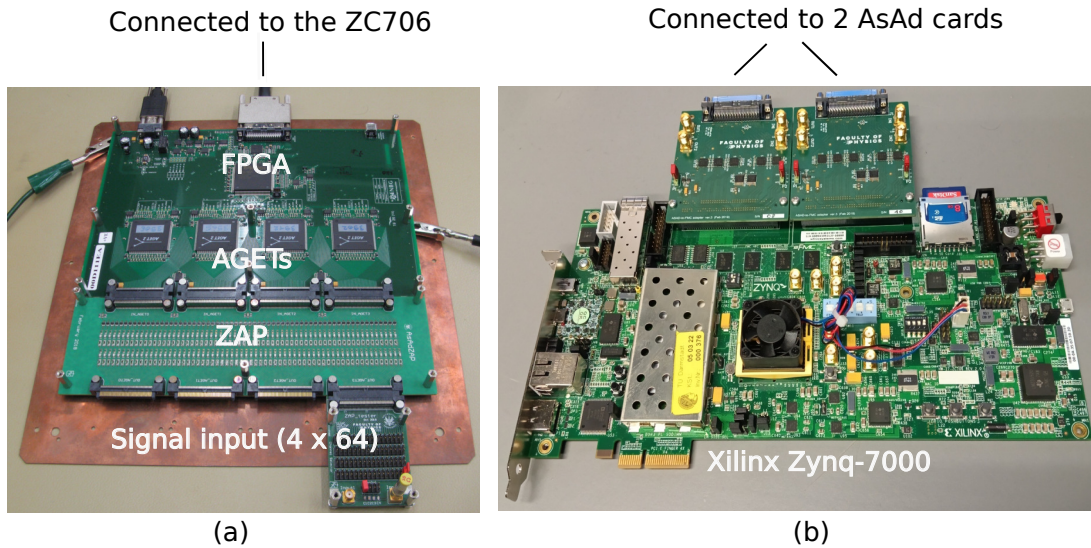


Figure 3.19.: (a) An AsAd board connected with a ZAP board (employed to protect the AsAd board). (b) The ZC706 evaluation board with the z-CoBo firmware [218] developed by Mikołaj Ćwiok and Marcin Zaremba from the University of Warsaw.

The core of the GET system is the front-end AGET chip, an ASIC (Application Specific Integrated Circuit) that incorporates 64 input channels. Analog signals from the readout anode are transmitted to these channels, each

comprising a charge-sensitive pre-amplifier (CSA), a shaper, an inverting $2\times$ gain amplifier, a discriminator, and a 512-cell switch capacitor array (SCA) in sequence.

The CSA integrates the detector's current signal, converting it into a voltage signal. It offers four gain settings, accommodating a wide dynamic range from 120 fC to 10 pC. The signal is subsequently shaped with a selectable peaking time between 70 ns and 1 μ s. The SCA, functioning as a buffer, continuously samples and stores the signal until an external trigger signal is received. The buffer duration for the signal is configurable. In addition to the input channels, four fixed-pattern noise (FPN) channels are included to extract intrinsic noise levels.

The analog signal, stored in the SCA, is digitized by the AsAd (ASIC and ADC) board with a 12-bit ADC. Each AsAd board integrates four AGET chips, as depicted in Figure 3.19(a).

The HYDRA-TPC employs a 1024-channel GET-based system, incorporating 4 AsAd cards. The digitized signals from the AsAd cards are transmitted to the concentration system based on the commercial Xilinx Zynq-7000 module implemented on the ZC706 evaluation boards, with the z-CoBo firmware [218], as displayed in Figure 3.19(b). This concentration system controls and coordinates the AsAd boards' operation and transfers data to external computers. The trigger rate is limited to a maximum of 1 kHz. The system's ENC performance was assessed by Yelei Sun at TU Darmstadt. The root-mean-square (RMS) noise was measured with and without a cable of 88.34 pF capacitance, as summarized in Table 3.1. The resulting ENC was found to be comparable with the AFTER ASIC, displayed in Figure 2.15.

	70 ns		1014 ns	
	RMS [ADC]	ENC [e^-]	RMS [ADC]	ENC [e^-]
Without cable	4.81 ± 0.29	975 ± 58	3.95 ± 0.99	792 ± 197
With cable	9.90 ± 1.27	1997 ± 254	8.81 ± 2.34	1764 ± 468

Table 3.1.: ENC measurements of the GET system with the peaking time of 70 ns and 1014 ns separately.

Multiplexing boards

Despite the implementation of the GET system, it remains insufficient for its 1024 channels to read signals from all 5632 pads. To overcome the limitation of the available number of channels, an intermediate multiplexing technique was employed.

A generic multiplexing [219] functions as dividing n pads into n/i groups, each containing i pads. A single electronic channel is shared among the n/i pads across groups. However, to identify the specific group containing the signal, i.e., to de-multiplex, an additional n/i channels are necessary. Therefore, the total number of electronic channels required is reduced from n to minimally $n/i + i$, significantly reducing the demands on the readout electronics. For the HYDRA-TPC, the pad plane is divided into eight segments, each comprising 32×44 pads, as illustrated in Figure 3.20 (Top). Multiplexing boards were constructed according to the scheme depicted in Figure 3.20 (Bottom). Each of these boards is directly inserted into the pad plane's pins and connected to two AGET chips via SAMTEC ERCD 1 m cables with a capacitance of 88.58 pF and a characteristic impedance of 50 Ω .

Each multiplexing board comprises 128 electronic channels. A group of four channels is associated with 22 pads in a half-column of the pad plane. The first pad within each half-column serves as a ground connection, allowing for electrical grounding to the chamber through a cable, as illustrated in Figure 3.16. The second pad is not multiplexed, serving as a reference for de-multiplexing. A detailed explanation of this process is provided in Section 6.2.3. The remaining 20 pads in the half-column are interconnected in groups of three, with each group connected to one of the four channels via conductive tracks.

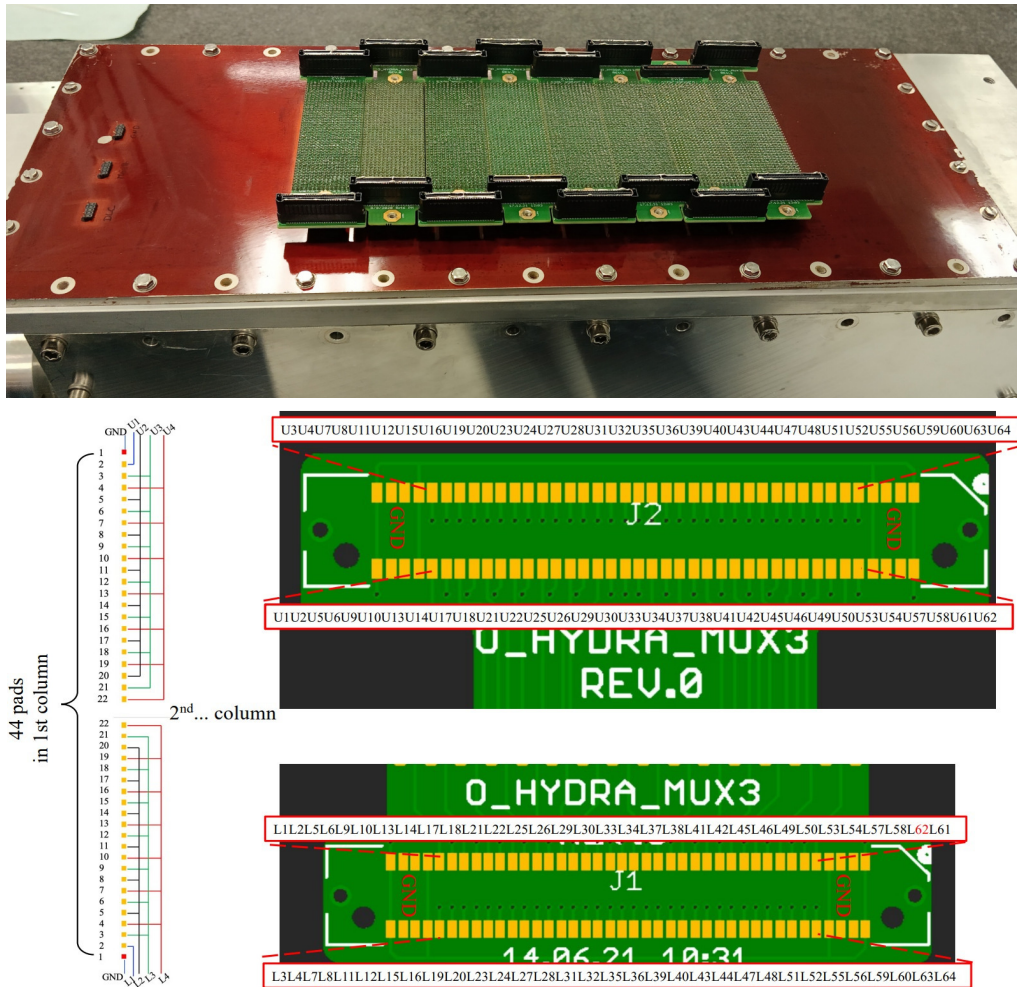


Figure 3.20.: (Top) Multiplexing boards installed on the pad plane. (Bottom) The scheme for one multiplexing board. The multiplexing boards are designed by Yelei Sun and constructed at the electronic workshop of TU Darmstadt.

4. Drift Field Simulation

To provide a guidance for the design of the HYDRA-TPC, two-dimensional simulations were conducted to evaluate the uniformity of the drift field within the active volume of HYDRA-TPC. This investigation utilizes the Finite Element Method and the Monte Carlo method. The active volume dimensions are referenced from Chapter 3. The detailed simulation of the amplification region, including the GEM and the resistive Micromegas, is currently not considered. The techniques of these detectors have been highly developed, therefore, their characteristics are directly adopted from manufacturer specifications and available literature.

4.1. Simulation Methods and Software

Finite Element Method

The Finite Element Method (FEM) is a widely-employed numerical technique for obtaining approximate solutions to problems involving complex structures [220]. FEM achieves this by discretizing the problem domain into smaller and simpler elements [221]. Within each individual element, algebraic equations are generated based on the element's size and properties. These element-level equations are then assembled into a comprehensive system of equations representing the entire domain with inter-element continuity. Finally, by incorporating boundary conditions, the equation of the global system is solved, and the results are presented in tabular or graphical formats.

Gmsh [222] is a powerful open-source finite element meshing tool for electric field simulation. It offers a user-friendly GUI interface and visualization with the feature of importing directly CAD drawings. It is also scriptable, which enables the customization of geometry. Notably, Gmsh offers a diverse range of meshing algorithms and a straightforward process for mesh refinement, catering to the specific needs of various simulations [223]. A customized geometry script usually contains the following as the input:

1. All geometrical entities, such as lines and surfaces, and physical groups, such as physical curves and surfaces. These definitions should be exclusive and comprehensive, ensuring that there is no overlap or connection between them and that no part of the entire space remains undefined.
2. Specifications of the meshing algorithm and its parameters. This crucial step dictates how the space is discretized into smaller elements inside different physical groups, directly impacting the accuracy and efficiency of simulations.

The meshing result by Gmsh can be directly imported to ElmerFEM [224] which is usually employed to calculate a wide range of physical processes for example fluid dynamics, heat dissipation, and electromagnetics based on the FEM. A user-defined input script provides ElmerFEM with a physical description of the meshing result generated by Gmsh. This script specifies the material properties and boundary conditions for each physical entity defined in Gmsh. Additionally, it designates the physical processes and equations to be employed in the calculations.

Monte Carlo Method and Garfield++

As discussed in Section 2.2.2, the drifting of electrons within a gas mixture under an electric field is a complex procedure involving a series of stochastic processes, such as collisions with gas molecules. These processes

often present significant challenges for analytical methods. The Monte Carlo (MC) method [225] offers an effective alternative by simulating random processes of electron drift. It is achieved by dividing the electron's entire drift path into small intervals. Within each interval, the random process is estimated and sampled from the probability distributions of the scattering process, the electron free path, and momentum change after collisions. This sampling process is controlled by the total collision rate at the electron energy of each interval. By iterating this process for a sufficiently large number of simulated electrons, the MC method builds a statistical model of electron drifting within the gas mixture.

Based on the MC method, Garfield++ [226], an open-source software package, specializes in the simulation of gaseous detectors and ionization phenomena. To simulate the motion of electrons within an electric field, the simulation setup using Garfield++ usually involves three primary components:

1. Definition of the geometry. The shapes, dimensions, and positions of all electrodes and insulators are defined.
2. Definition of material properties. The material composition for each insulator of interest, including the gaseous mixture within the active volume, is defined, along with its associated properties such as density, pressure, relative dielectric constant, and temperature.
3. Calculation of the electric field. Boundary conditions, including potential values on electrodes, are applied. The electric field is calculated in the active volume.
4. Electron source and physics processes. Initial conditions are defined for the simulated electrons, including their starting energy and spatial distribution. Additionally, a transport model is specified, incorporating relevant cross-sections for physical processes such as elastic and inelastic collisions between electrons and gas molecules.

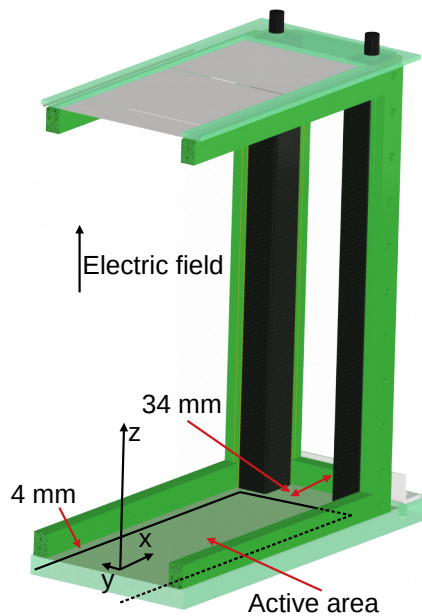


Figure 4.1.: Sectional view of the field cage (schematic). Axis z denotes the direction of the drift field. The black lines on the cathode denote the active area.

4.2. Drift Field Simulation

To ensure the uniformity of the drift field within the active volume of the HYDRA-TPC, a two-layer wire configuration is employed. Considering the proximity of the wires to the grounded chamber walls, particularly for those near the cathode (approximately 20 mm) where the potential difference is maximum, a single layer of wires leads to field line distortion between neighboring wires. The utilization of two wire layers can potentially mitigate this effect.

As illustrated in the cross-sectional Figure 4.1, the active volume of the HYDRA-TPC is positioned 34 mm from the narrow sides of the field cage. Given this distance, edge effects on the electric field homogeneity are negligible at these sides. In contrast, the active volume is situated merely 4 mm from the inner layer of wires on the wider sides facing the TPC windows. In this configuration, edge effects cannot be disregarded. As maintaining electric field homogeneity is essential for accurate tracking performance, minimizing field distortions near these wires became a primary consideration for the simulation.

Moreover, generating a sufficiently dense mesh for a full three-dimensional model for the field cage can significantly increase computational costs during the meshing process itself. Consequently, a two-dimensional analysis of the field uniformity across a central cross-section of the TPC was undertaken. This approach simplifies the simulation while focusing on the essential challenge of the drift field's uniformity.

4.2.1. Two-Dimensional Meshing

As outlined in Section 4.1, the geometry for the two-dimensional simulation incorporates definitions of entities and their boundaries for electrodes and distinct volumes.

The outermost layer of rectangular and polygonal shapes, representing the chamber and metal core pad plane, defines the 'world'. The cathode and its PCB are positioned above the chamber's bottom plate. Wires (represented as small closed circles in Figure 4.2(a), simplified as black dots in the center) and strips (represented by mini rectangles) are vertically aligned above the cathode, with an interval of 3 mm. The GEM (a thin rectangle above the top strips) and the Micromegas (beneath the metal core pad plane) are located between the top strips and the metal core pad plane. The whole geometry was carefully examined to preclude any open or overlapping shapes. As visualized in Figure 4.2(a) by Gmsh, the simulation considers the realistic geometry, although the potential bending of window foils under overpressure inside the chamber is not accounted for. Furthermore, in this section, the GEM and Micromegas are modeled as flat conductive planes, neglecting their micro-structures.

The meshing result is depicted in Figure 4.2(b). The calculation employs the triangular meshing method. A color map, ranging from red to green, visually represents the mesh quality within the elements. By default, red represents a low-quality mesh, while green indicates a high-quality mesh. Within the region of interest (mostly the active volume), the mesh exhibits consistently high quality (green). Although some peripheral regions, such as the PCB for the cathode, possess a moderate mesh quality (yellow or brown), it barely influences the subsequent electric field calculation.

As a result, there are regions where the triangles are less dense. The approach, where a coarser mesh is utilized in less critical areas (for instance, the middle part of the field cage), aims to achieve a more than 50% reduction in computational time. This optimization is justified because the expected variations in the electric field within these regions are less pronounced than in the regions where it is close to any electrode or dielectric boundaries. However, a reduction in mesh density throughout the remaining regions is not feasible, particularly in close proximity to the wires and strips. The finite size of the electrodes excludes the utilization of larger mesh elements. Additionally, the significant variations in the electric field within this vicinity require the employment of a fine mesh structure.

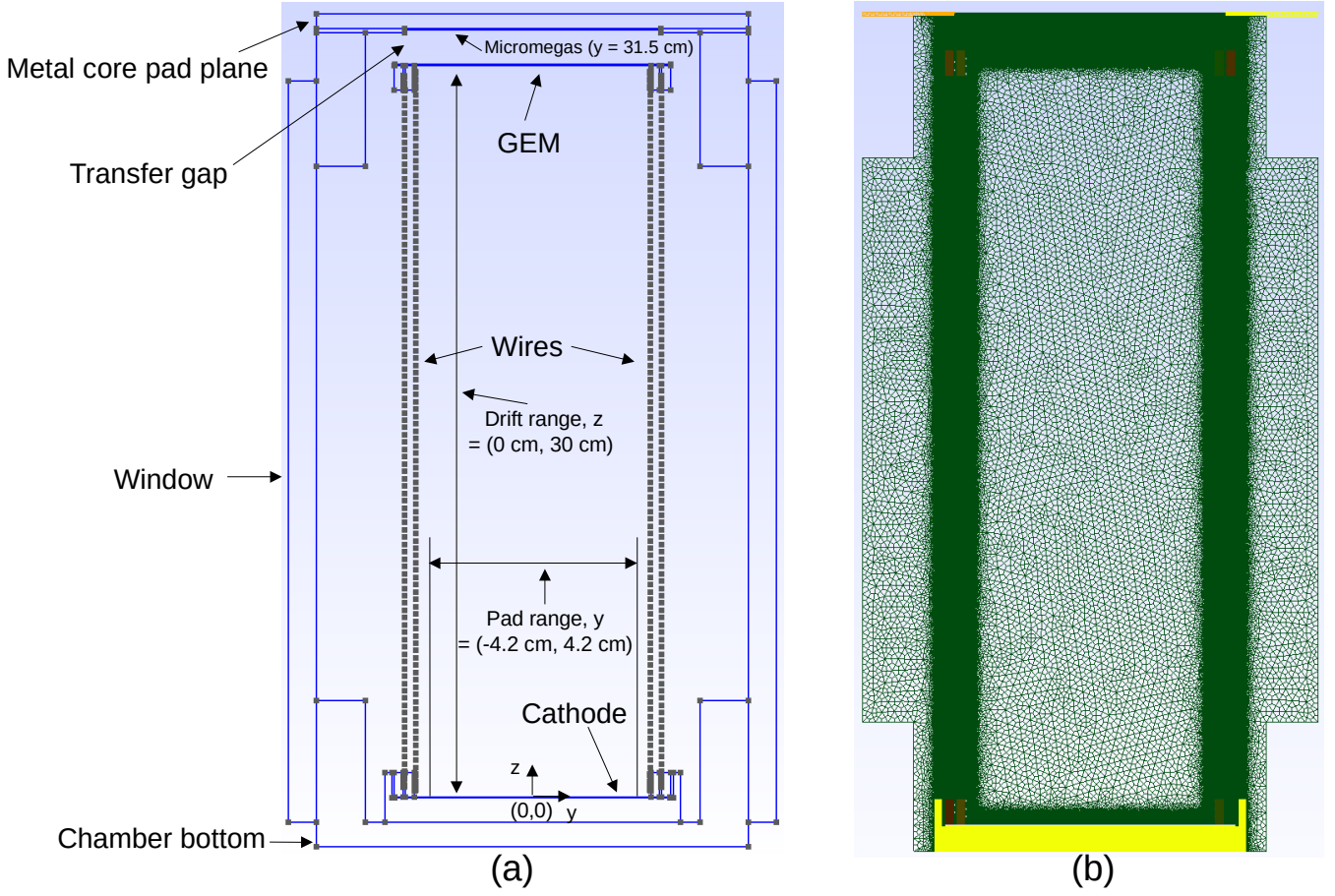


Figure 4.2.: (a) Two-dimensional geometry input file for Gmsh and (b) result with the triangular meshing method. More details are in the text.

4.2.2. Drift Field Calculation

The meshing results were subsequently imported into ElmerFEM to calculate the electric field and the potential. The calculation is based on the electrostatic model [227], considering solely the electric part of Maxwell's equations,

$$\nabla \cdot \vec{D} = \rho, \quad (4.1)$$

$$\nabla \times \vec{E} = 0. \quad (4.2)$$

The relative permittivity ϵ of the gas is set to 1.00058, and that of epoxy resin is set to 4.5. From Equation 4.1 and the electric field in the stationary scenario $\vec{E} = -\nabla\phi$, the equation for the electrostatic calculation is given by

$$-\nabla \cdot \epsilon \nabla \phi = \rho. \quad (4.3)$$

Dirichlet boundary conditions are applied to the calculation in which the potential $\phi(r_i)$ of the electrode surface stays constant. The chamber, including the window and the pad plane, is grounded. The distribution

of high voltages on electrodes is shown as follows:

Cathode — 7400 V,
GEM_{top} — 800 V,
GEM_{bottom} — 420 V,
Micromegas — 0 V.

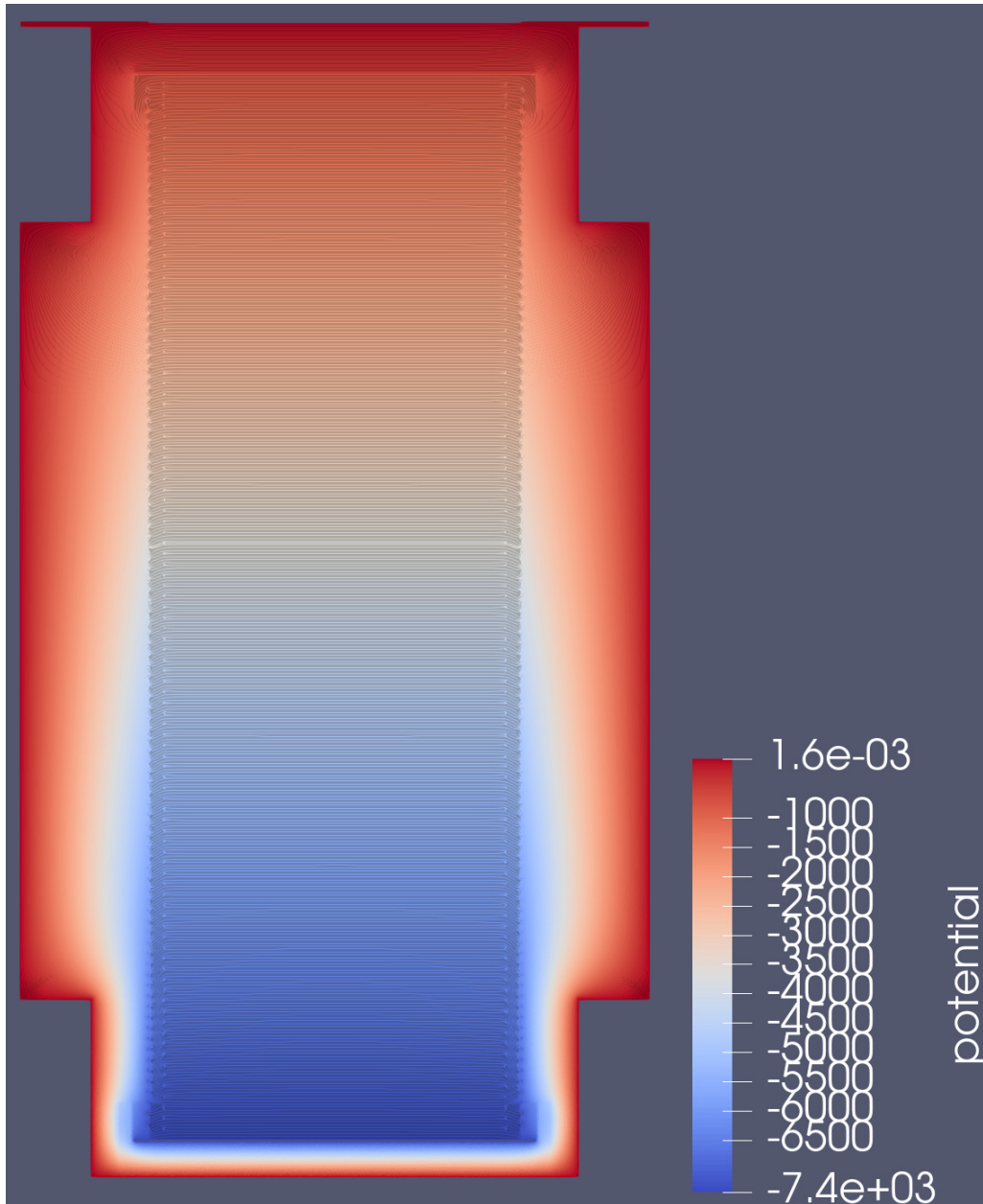


Figure 4.3.: Equipotential lines calculated from ElmerFEM (Unit: V).

A total of 394 wires and strips, every two of them being paralleled, are differentially biased with high voltage based on their distance from the cathode. A potential difference of 33 V exists between neighboring wires or

strips, leading to a potential gradient of 220 V/cm in the active volume. The potential difference between the last strips and the GEM is calculated to maintain consistency with the potential gradient. Visualization of equipotential lines, generated by ElmerFEM, is presented in Figure 4.3. In the central region of the active volume, the electric field is perfectly uniform, while, notably, due to the drastic potential difference between the chamber wall and the wires, there are significant field distortions between two layers of wires, which justifies the necessity of the employment of two layers of wires.

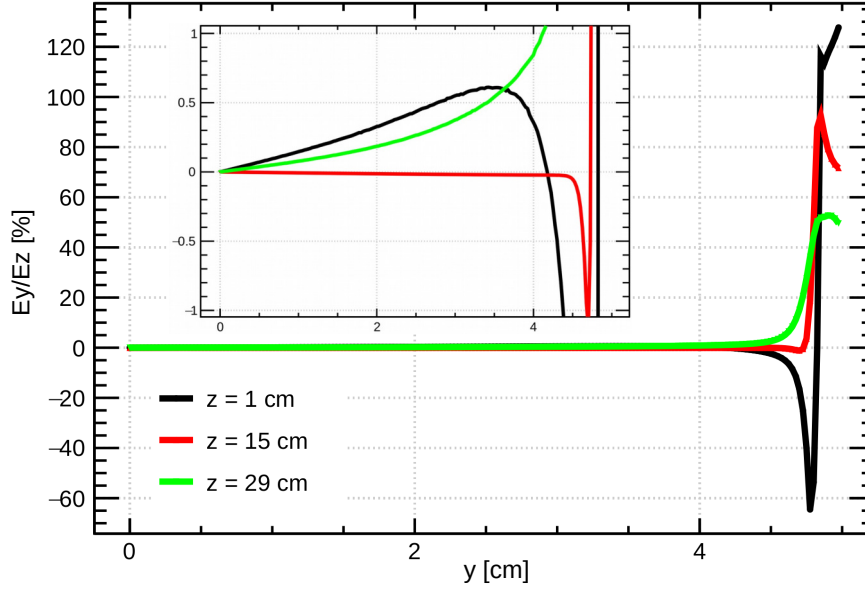


Figure 4.4.: Ratio E_y/E_z at various positions. The horizontal position y ranges from 0 (at the center of the cathode) to 4.2 cm (at the edge of the active area). The vertical positions z are in the center of the drift field, 1 cm from the cathode and from the GEM respectively.

To quantify the drift field homogeneity, the ratio of the horizontal component (paralleled to the cathode, E_y) to the vertical component (perpendicular to the cathode, E_z) of the electric field is calculated at different horizontal (y) and vertical (z) positions of the field cage as presented in Figure 4.4. A high absolute E_y to E_z ratio indicates a large field distortion, while zero suggests uniformity. There are three regions:

1. Inside the active volume ($y < 4.2$ cm). The distortion is negligible where it is at the center of the active volume ($y = 0$). The field remains uniform for the red line until $y = 4.2$ cm, while the ratio rises gradually for black and green curves. At the margin of the active volume where $y = 4.2$ cm, the ratio reaches 0.8%.
2. Between the active volume and the inner layer of wires ($4.2 < y < 4.8$ cm). A deviation from field uniformity is observed for all z positions. This is primarily due to the influence of the wire and strip geometry. Furthermore, near the GEM and cathode planes, local field distortions can occur due to dielectric discontinuities between the gas mixture and the epoxy resin of the supporting PCB frames.
3. Between two layers of wires ($4.8 < y < 5.3$ cm). The ratio experiences a significant increase to some maximum value and decreases at a larger horizontal position. This corresponds to the compression of the equipotential lines around the outer layer of wires in Figure 4.3. Locally the potential reaches the extremum at the closest to the wires.

4.2.3. Drift field with different resistor chains

In Section 4.2.2, electrodes are assigned high voltage with nominal values. However, in reality, the construction of resistors is inherently subject to uncertainties. Due to manufacturing processes, it is only possible to constrain the maximum deviation from the nominal value, known as tolerance. This tolerance arises from various factors, including the shaping of the resistors and material impurities. Manufacturers typically offer resistors with tolerances ranging from $\pm 0.01\%$ to $\pm 20\%$.

To assess the impact of various resistance tolerances (0.1%, 1%, 3%, 5%, 10%) on field homogeneity, a simulation of the drift field was conducted with different series of resistors. In each series, resistor values were randomly sampled from a truncated Gaussian distribution. The mean of the distribution was set to the nominal resistance value, while the standard deviation was set to half the tolerance value to ensure that approximately 95% of the generated values fell within a two- σ range. The distribution was truncated to limit the maximum deviation to the specified tolerance.

In this and the following simulations, the pseudo-random generator seed is calculated by the system time at execution to ensure the independence of generated values on the seed and break the inner correlation among different simulations.

Figure 4.5 displays a ratio E_y/E_z of more than 1.7% for the electric field with resistance tolerance equal to or larger than 5%. While the tolerance is less than 5%, the ratio E_y/E_z does not vary significantly, even between tolerances of 0.1% and 3%. This can be attributed to the nature of the Gaussian distribution, which concentrates the majority of samples near the mean value, thereby mitigating the impact of tolerance differences.

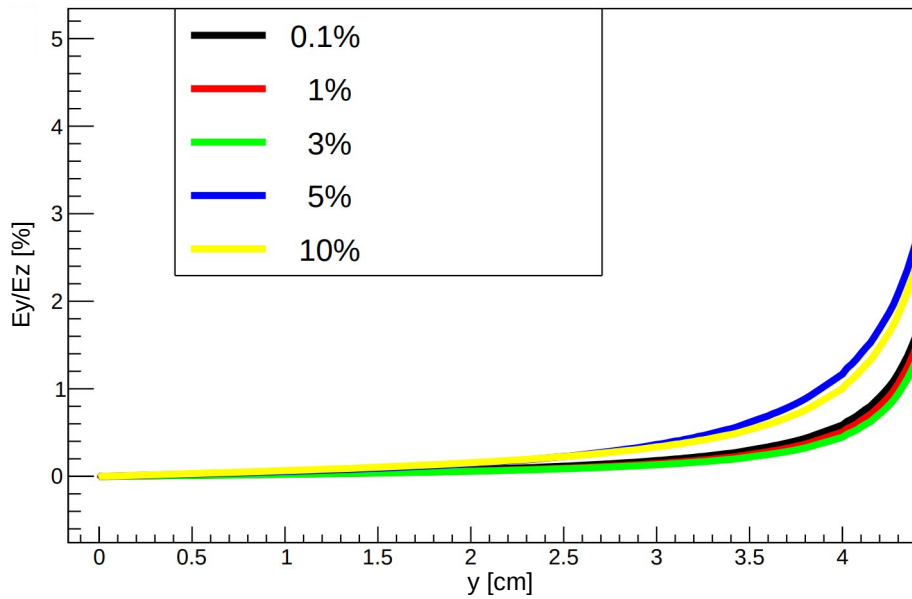


Figure 4.5.: Ratio E_y/E_z for various construction tolerance of resistors. The horizontal position y ranges from 0 (at the center of the cathode) to 4.2 cm (at the edge of the active volume). The vertical position is fixed to $z = 1$ cm from the GEM.

4.2.4. Calculation of the Transfer Field

The transfer field is responsible for the extraction and transport of amplified electrons from the GEM towards the Micromegas; therefore, the transfer gap should be minimized to ensure minimal position deviation of the

electrons upon their drift-through. On the other hand, the minimum distance is constrained by the geometry, including the incompressible thickness of the GEM-PCB.

Computational simulations have been conducted to analyze the transfer field under various GEM-to-Micromegas distances. To simplify the calculation, both the GEM and the Micromegas were considered smooth, conductive plates. For a distance of 5 mm, field distortion is almost negligible, as illustrated in Figure 4.6(a). The ratio E_y/E_z is at a maximum of 1.2% at the edge of the active area and the proximity of the boundary between the PCB and the gas volume where the field is distorted due to the discontinuity of the dielectric constant. In contrast, a GEM-to-Micromegas distance of 11 mm exhibits a maximum distortion of 4.9%.

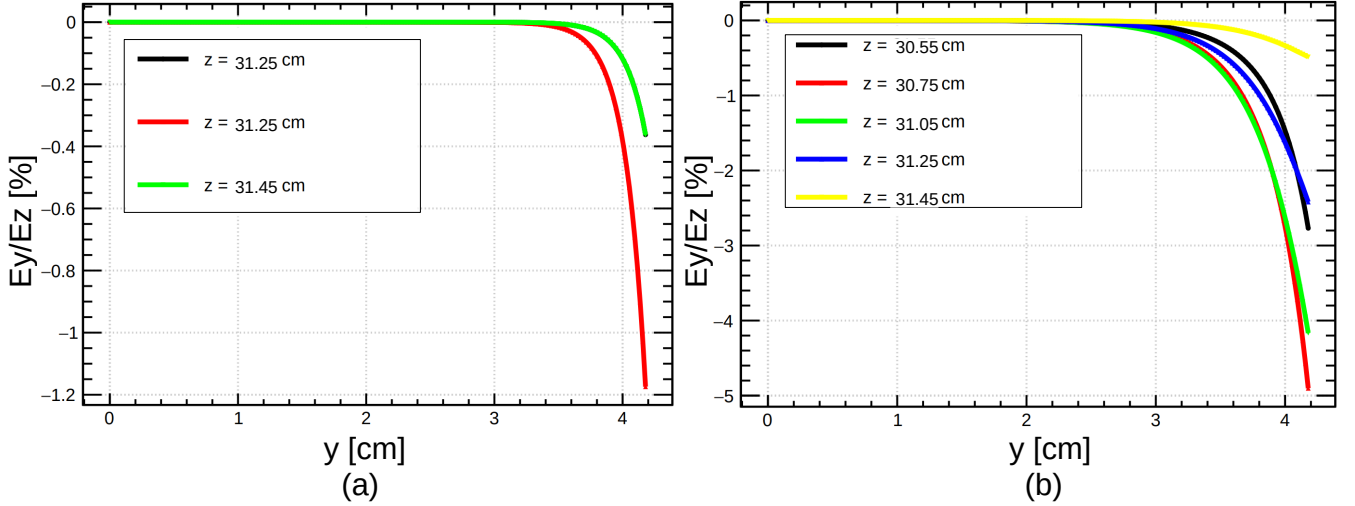


Figure 4.6.: Ratio E_y/E_z for a GEM-to-Micromegas distance of (a) 5 mm, with the bottom surface of the GEM located at $z = 32.00$ cm. and (b) 11 mm, with the bottom surface of the GEM located at $z = 31.40$ cm. The horizontal position y ranges from 0 (at the center of the GEM) to 4.2 cm (at the edge of the active volume). z denotes the distance from the cathode, as showcased in Figure 4.2.

4.2.5. Discussion

The simulation results indicate that the drift field uniformity within the active volume, considering the geometry depicted in Figure 4.2, is limited to a variation of less than 1%. For the construction of the flexible PCBs, resistors with a 1% tolerance were selected for the resistor chain. This choice was based on the observation that utilizing resistors with a 0.1% tolerance does not significantly improve field uniformity, as illustrated in Figure 4.5. Furthermore, resistors with a 1% tolerance are the most readily available and economical among all tolerance options. In reality, the drift field uniformity can be degraded by various factors, including space charge effects in the drift volume and mechanical deformations such as wire bending. Calibration measurements utilizing laser tracks are highly desirable for accurately characterizing the drift field.

During the assembly of the HYDRA-TPC, several deviations from the geometry in Figure 4.2 were introduced. The assembly of the cathode, entire field cage, and the GEM were raised and supported by spacers, reducing the GEM-to-Micromegas distance to 4.2 mm. This modification is not expected to impact the conclusions of this section, as the field cage is further from the chamber's bottom plate, which is grounded, potentially improving field uniformity. Additionally, the simulation model includes epoxy resin around the strips near the

cathode to enhance field uniformity, a feature not implemented in the physical construction. This discrepancy does not affect the results of this thesis, as neither charged particles of interest nor laser beams can be directed to this region, thereby not influencing the drift of primary electrons.

4.3. Electron Drift Simulation

To further evaluate the adequacy of the less than 1% field distortion obtained in Section 4.2.2, three-dimensional simulations of electron drift within the active volume were conducted using Garfield++.

The HYDRA-TPC is supposed to serve under the magnetic field. However, in this section, only the electric field was considered in the drift simulation. The primary objective was to assess the impact of electric field non-homogeneity on electron drift. The effects of non-uniform magnetic fields and perpendicular magnetic field components will be investigated in a separate and dedicated simulation, which is beyond the scope of this thesis.

4.3.1. Simulation Setup

The two-dimensional electric field distribution within the active volume, as calculated by ElmerFEM, can be incorporated into the simulation using the ComponentElmer2d class of Garfield++. The three-dimensional electric field is thus generated by replicating the two-dimensional field distribution along the third dimension. The class also defines the dielectric properties of various volumes and the gas composition within the active volume.

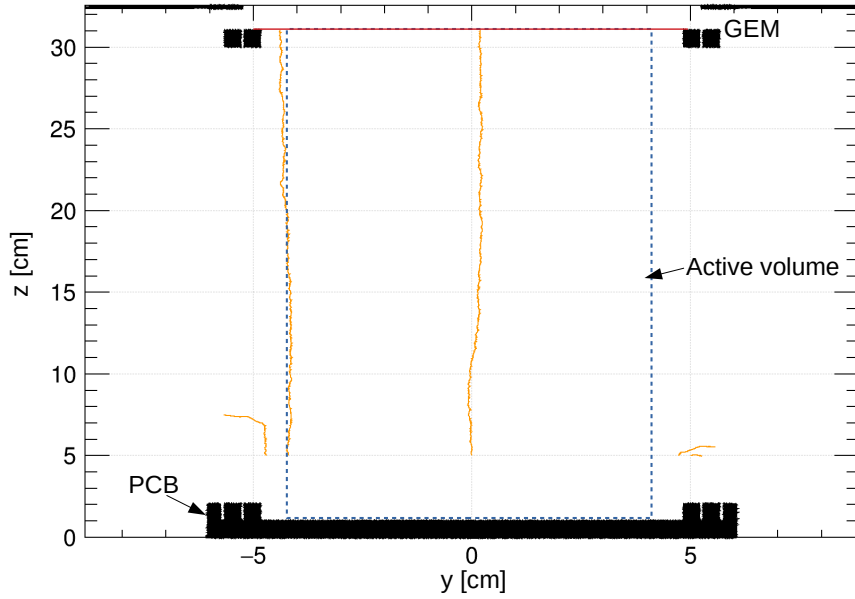


Figure 4.7.: The tracks of drift electrons starting at different y positions (from left to right, one at 0.5 cm from the inner layer of wires, one at the edge of the active volume, one inside the active area, two between the inner layer and the outer layer of wires). These electrons are emitted at $z = 5$ cm and are absorbed by the GEM or the wires. The black area is epoxy resin.

The AvalancheMicroscopic class tracks the trajectory of each drifting electron on a collision-by-collision basis. The electron transport parameters, such as scattering cross-sections for elastic and inelastic collisions between

electrons and gas atoms, as well as electron exchange interactions, are calculated using the Magboltz class [228]. These parameters are crucial for accurately simulating the electron's interaction with the gas medium. Several instances of electron tracks are displayed in Figure 4.7.

4.3.2. Position Deviation of Drift Electrons

Following their trajectories through the active volume, a discrepancy can arise between the initial emission position y_i and the final arrival positions y_f of the drift electrons. This deviation can be attributed to two primary factors:

1. Diffusion. As detailed in Section 2.2.2, the random scattering of electrons by gas molecules induces a transverse motion during their drift towards the GEM.
2. Displacement. Drift trajectories of electrons follow electric field lines. Consequently, even minor distortions in the field, such as an E_y/E_z ratio of 1%, can lead to a collective deviation $y_i - y_f$ for all electrons.

Figure 4.8 presents the distribution of positional deviations extracted from a drift simulation of 10000 electrons, utilizing the field configuration calculated in Section 4.2.3 based on resistors with a 1% tolerance. This field configuration is considered the default drift field throughout this section. The gas mixture employed is the standard T2K composition (95% Ar + 2% iso-C₄H₁₀ 3% CF₄). The displacement of the electrons is quantified by the mean value of the distribution, 0.35 ± 0.02 mm, while the standard deviation, 1.68 ± 0.01 mm, represents the diffusion.

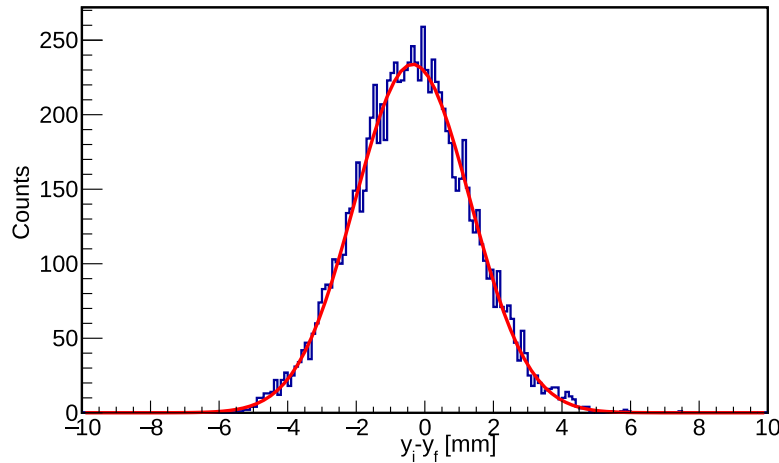


Figure 4.8.: Distribution of Position deviation for 10000 electrons drifting 3 cm from the cathode at the edge of the active area.

Both displacement and diffusion can impact the tracking resolution of the HYDRA-TPC. Moderate diffusion can contribute to the precise determination of the charge centroid and enhance spatial resolution as detailed in Section 2.3.3. Displacement, on the other hand, directly hinders precise tracking as the deviation displaces the charge centroid.

The deviation depends primarily on the properties of the working gas, as described in Section 2.2.2, such as pressure, temperature, and mixture components, and the uniformity of the electric field. Simulations of electron drift through the field cage were conducted using two gas mixtures, namely the T2K gas and the gas mixture of 90% Ar + 10% CO₂, as depicted in Figure 4.9. In these simulations, all electrons were initialized 4

cm from the cathode, the closest distance within which charged particles are not expected to penetrate, given the TPC's acceptance. The results indicate that the gas, 90% Ar + 10% CO₂, generally exhibits both smaller displacement and diffusion compared to the T2K gas. Notably, both gases demonstrate displacements below 200 μm across most positions within the field cage. Only at the extreme edge of the field cage ($y = 4.2$ cm) does the displacement reach 205 ± 17 μm .

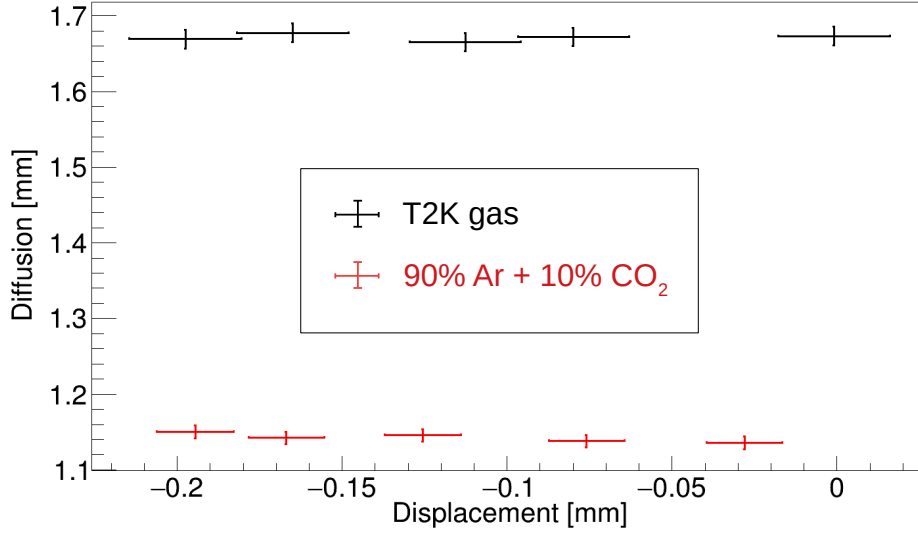


Figure 4.9.: Positional deviation of electrons in the drift region for different gas mixtures. The data points for each gas mixture, correspond to distinct y positions within the field cage (From left to right, $y = 4.2, 3.2, 2.2, 1.2$, and 0.2 cm). The default electric field is used. For both gas mixtures, the temperature is 293 K and the pressure is 1 bar.

4.3.3. Drift Deviation in the Transfer Field

Similar simulations were performed for the transfer field to quantify the positional deviation of electrons within this region. The electric field distributions calculated in Section 4.2.4 for both 5-mm and 11-mm GEM-to-Micromegas distance (d_t) were utilized to investigate electron drift behavior under different transfer fields. For this purpose, electrons were assumed to be emitted from various y positions on the GEM surface. To simplify the initial conditions, it was assumed that electrons were emitted with negligible initial velocity and zero spatial dispersion. This simplification allows for a focused evaluation of electron drift in the transfer field. It should be noted that in reality, electrons extracted from the holes of the GEM possess a distribution of initial velocities and angles.

Simulation results indicate that a 5-mm d_t can enhance electron extraction performance within the transfer field, as illustrated in Figure 4.10. Consistent with previous simulations for the drift field, the gas, 90% Ar + 10% CO₂, exhibits lower diffusion compared to the T2K gas. For $d_t = 5$ mm, both gas mixtures demonstrate negligible displacement in the central active region, while only minor displacement of less than 50 μm is observed at the edge.

In contrast, a larger distance of $d_t = 11$ mm leads to increased diffusion for both gases. Furthermore, at larger y positions (3.2 and 4.2 cm), a significant displacement of approximately 350 μm is observed, presenting a poorer extraction performance.

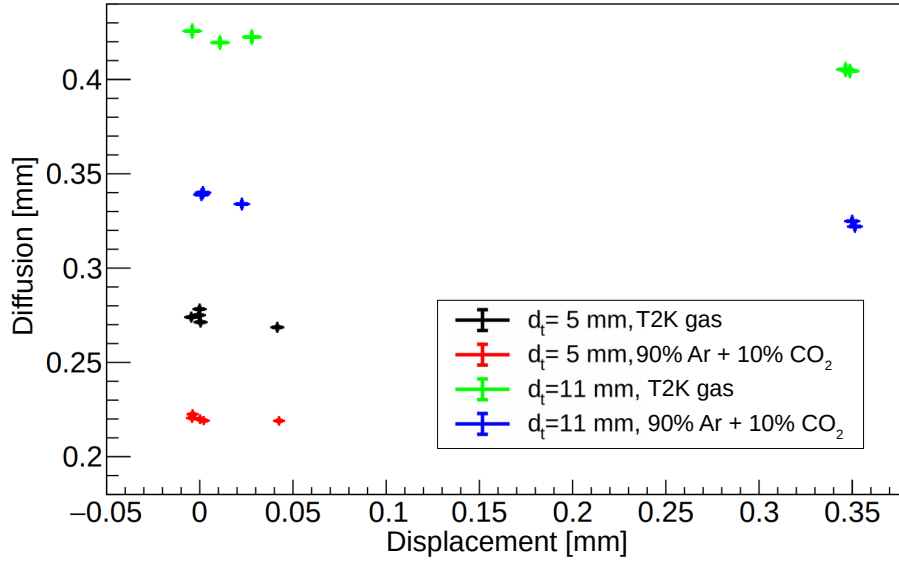


Figure 4.10.: Positional deviation of electrons for different d_t with the same electric field strength 1000 V/cm. For each point, there are 10000 electrons released at different y positions (From the left to right, $y = 0.2, 1.2, 2.2, 3.2$, and 4.2 cm).

4.3.4. Discussion

Combining simulation results from both the drift field and the transfer field, the total charge diffusion on the pad plane was estimated for GEM-to-Micromegas distances of 5 mm and 11 mm. A 50 μm diffusion contribution from each amplification stage of the GEM and Micromegas was considered. For the T2K gas, the estimated total charge diffusion is 3.19 mm and 3.20 mm, respectively, while for the 90% Ar + 10% CO₂ gas mixture, it is 2.94 mm and 2.95 mm, respectively. This charge diffusion on the pad plane corresponds to approximately two pads in the HYDRA-TPC, which is ideal for precise centroid calculation.

However, the impact of electron displacement cannot be disregarded. While the displacement in the drift region is negative, indicating that electrons are displaced further from the active region, the displacement in the transfer region is positive, suggesting an inward displacement. These results are consistent with the calculated field distributions in both regions. The overall displacement, considering both regions, reaches a maximum of approximately 210 μm for the gas, 90% Ar + 10% CO₂, and 230 μm for the T2K gas at the edge of the active region. For the majority of the active volume, the spatial resolution resulting from electron deflection is less than the requirement of 200 μm in Section 3.1.2.

To further enhance spatial resolution, optimization of the working gas for the HYDRA-TPC is essential, as evidenced by the better performance of the gas, 90% Ar + 10% CO₂, compared to the T2K gas in terms of both electron diffusion and displacement. Additionally, a field correction map, derived from laser reference measurements, would be beneficial, particularly for the edge regions.

5. Characterization of HYDRA-TPC

This chapter presents a detailed performance evaluation of the HYDRA-TPC. The primary emphasis is on characterizing the amplification properties of the TPC. This is a realistic assessment of the overall amplification of the integrated detector of the GEM and the Micromegas. While the uniformity of the drift field was assumed based on simulations presented in Chapter 4, a direct experimental verification was not conducted within the scope of this thesis.

5.1. Test Preparation

5.1.1. Test Setup

Characterization tests were conducted in the detector laboratory at GSI. The HYDRA-TPC was continuously flushed with a test gas mixture composed of 90% argon and 10% carbon dioxide. The gas flow rate was maintained at approximately 12 L/min to minimize impurity levels within the chamber during initial setup. During stable operation, the flow rate was reduced to 6 L/min to ensure slightly above-atmospheric pressure and to mitigate oxygen contamination arising from leakage and outgassing of materials inside the TPC. The temperature was kept constant at approximately 22 degrees Celsius. Oxygen levels were monitored and consistently remained below 30 ppm.

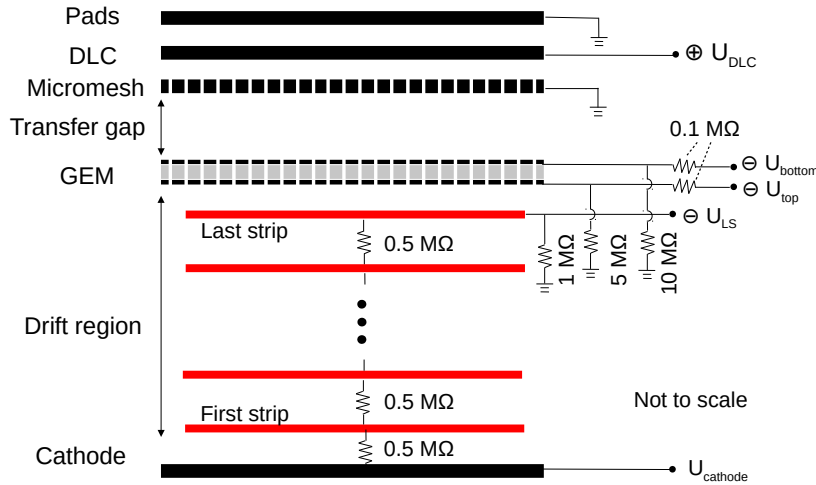


Figure 5.1.: Scheme for electrical connection and resistors of the HYDRA-TPC for gain tests. Red rectangles are wires and strips.

High voltage (HV) was subsequently applied to the TPC. As illustrated in Figure 5.1, five electrodes were connected to HV inputs, with the cathode voltage reaching a maximum of 8 kV. Electrodes, namely the cathode (U_{cathode}), the last strip (U_{ls}), the top side of the GEM (U_{top}), and the bottom side of the GEM (U_{bottom}), were

connected to a 4-channel 8-kV power supply, CAEN N1470ET, while the DLC (U_{DLC}) was connected to a separate CAEN N471A power supply. To ensure safe charge evacuation in the event of a power supply trip, resistors of different values were connected to the ground in parallel for both sides of the GEM and the last strip, as showcased in Figure 5.1. The resistor selection was based on two primary considerations:

1. Preventing sparks. RC time constants were optimized to ensure uniform discharge of both sides of the GEM upon a trip. This balanced discharge minimizes the voltage difference across the GEM, thereby reducing the risk of sparks.
2. Current direction. Resistor values must be chosen to avoid reverse current flow into the power supply. Power supplies are designed to provide instead of sinking current. For the last strip, the flow of current to the power supply (I_1) is given by

$$I_1 = I_0 - I_2 = \frac{U_{\text{ls}}}{R_{\text{ls}}} - \frac{U_{\text{cathode}} - U_{\text{ls}}}{R_{\text{FC}}}, \quad (5.1)$$

where I_0 and I_2 denote the currents through the resistor R_{ls} and the field cage with a total resistance $R_{\text{FC}} = 24.75 \text{ M}\Omega$. To ensure a positive value for I_1 , the first term on the right-hand side of Equation 5.1 is required to be greater than the second term.

To mitigate the risk of discharge between electrodes, a controlled ramping rate was implemented during the power-up and power-down steps of the power supplies. The DLC power supply ramping rate was typically limited to 50 V/s while both sides of the GEM were ramped at a rate below 30 V/s. Given the proximity of the last field cage strip to the GEM's top side with a distance of 2.5 mm, a synchronized ramping of the last strip and cathode, along with the GEM, is essential to prevent significant localized potential differences.

Furthermore, a trip connection was established across all four channels of the N14700ET power supply. In the event of electrode discharge, resulting in an excessive current above a predefined threshold, the entire power supply is immediately shut down. The predefined upper limits for the currents can be calculated with resistors in Figure 5.1.

Signals were processed and digitized using a single-channel electronics system consisting of an ORTEC 421 pre-amplifier, a Tennelec TC243 amplifier, and an ORTEC 928 MCB ADC module. The digitized data were subsequently transferred to a personal computer for further analysis. The ADC spectrum generated by the module was visualized using Maestro software.

5.1.2. Current Measurements

The DLC layer of the Micromegas and the GEM were connected to a pico-ammeter which monitors the currents on these electrodes. During standard operation for the HYDRA-TPC, all electrodes within the field cage and the GEM are always maintained at a negative bias potential, while the DLC layer is positively biased. The specific high voltage value applied to each electrode is determined based on the desired electric field strength within each region of the TPC, as exemplified in the following configuration:

$$\begin{aligned} E_d &= 220 \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 350 \text{ V}, \\ E_t &= 1000 \text{ V/cm}, \\ U_{\text{DLC}} &= 440 \text{ V}, \end{aligned}$$

where E_d represents the electric field in the drift region, ΔU_{GEM} representing the potential difference between the top and bottom sides of the GEM, E_t the transfer field and U_{DLC} and the potential bias on the DLC layer.

Upon application of a high voltage, the pico-ammeter continuously monitors the current on the DLC layer and displays the one-second averaged value. During the voltage ramp-up process, the current typically quickly increases to a maximum of approximately 200 nA. Subsequently, as the voltage reaches the target value, the current gradually decreases to a stable value which is referred to as the leakage current. As depicted in Figure 5.2, one measurement showcased that the current spiked to ~ 180 nA and then decreased to ~ 0.9 nA where it remained stable. The magnitude of the leakage current exhibits a dependence on the electric field strength within the amplification region [229]. Furthermore, high humidity levels and contamination, such as dust, have been shown to increase leakage currents significantly [230].

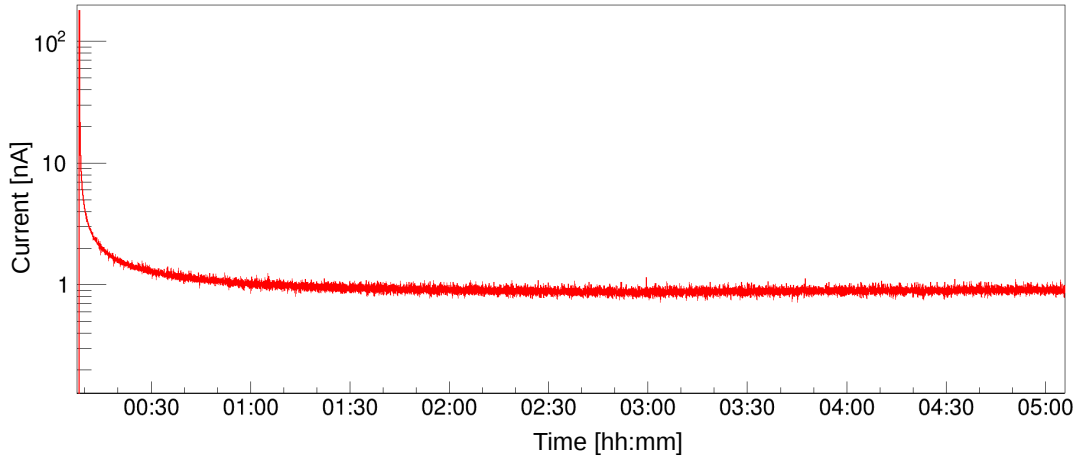


Figure 5.2.: Measurement of the leakage current on the DLC with a bias of 440 V.

5.1.3. ADC Spectra of ^{55}Fe Source

A ^{55}Fe source was employed to perform the test measurements. The corresponding ADC spectra are presented and analyzed in this section.

^{55}Fe X-ray Source

The ^{55}Fe source has a half-life of 2.7 years and decays into ^{55}Mn by K-shell electron capture. Subsequently, a higher shell electron fills the vacancy mainly by emitting Auger electrons with a probability of 60% and emitting X-rays (25% are 5.9 keV K_α X-rays and 3.4% are 6.5 keV K_β X-rays) with a probability of 28% [231]. During the source tests of the TPC, KX-rays are the desired channel.

The source was positioned external to the chamber, directly facing the detector window with a distance of 20 cm as depicted in Figure 5.3. It was aligned to the vertical and horizontal midpoint of the active volume of the TPC.

ADC Spectra of the X-ray Signals

The Tennelec TC243 amplifier provides two adjustable parameters: shaping time and amplification gain. The shaping time was set to the amplifier's minimum value of $2\ \mu\text{s}$ to prioritize fast signal processing and minimize dead time between events. The available gain options are $5\times$, $10\times$, $20\times$, and $50\times$. An amplification gain of $20\times$ was determined optimal through a series of tests. Lower gains ($5\times$ and $10\times$) resulted in inadequate signal amplification, obscuring the signal within the background noise observed in the ADC spectrum. Conversely, a gain of $50\times$ led to significant electronic noise, degrading the signal.

Figure 5.4 illustrates an ADC spectrum of X-rays on the micromesh. Apart from the noise, there is only one peak which corresponds to K_α X-rays as it is the most pronounced peak. Potentially additional peaks may be

observed in the ADC spectrum of ~ 6 keV X-rays. These include the photon escape peak, arising from the escape of lower-energy X-rays generated during the photoelectric effect, the K_{β} X-ray peak, and the summation peak, resulting from the detection of multiple photons by coincidence. However, these peaks are indistinguishable due to the dominance of the K_{α} X-ray peak and its large diffusion.

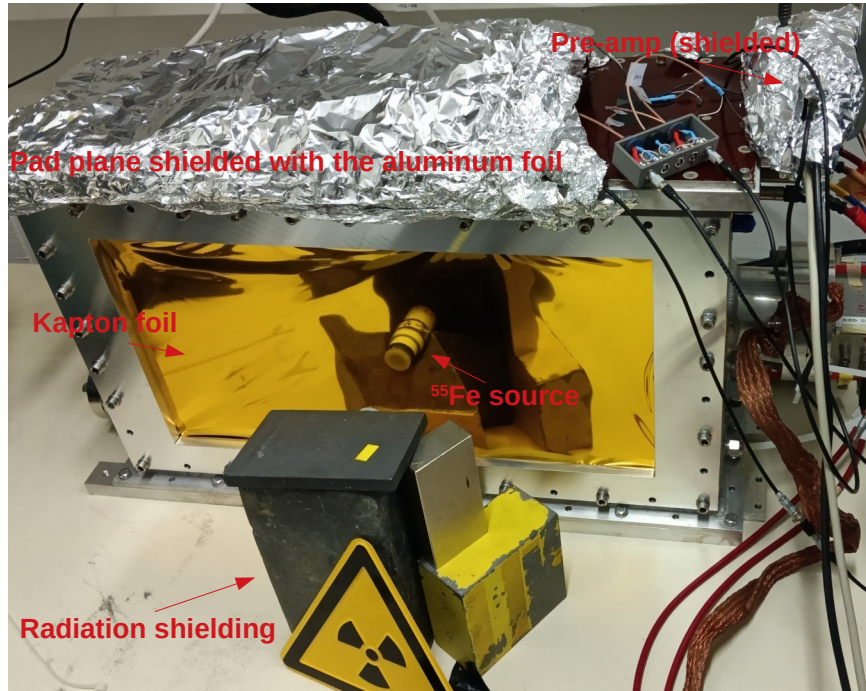


Figure 5.3.: Setup for the gain tests with the ^{55}Fe source.

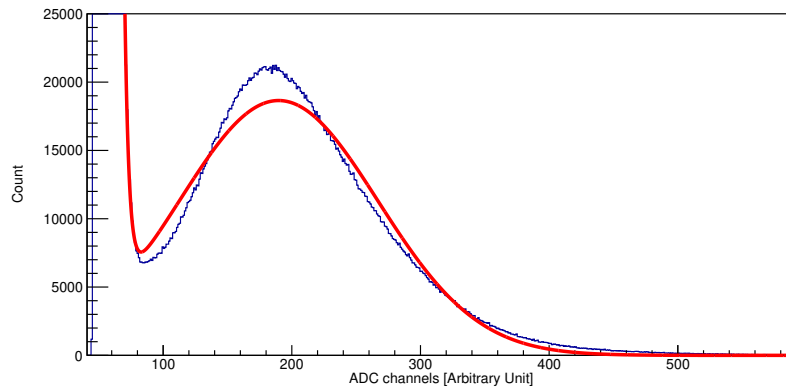


Figure 5.4.: X-ray ADC spectrum on the micromesh. The test duration is 181 s. The electronics gain is $4\times$. The blue line is the raw spectrum while the red curve is the fitting result.

A function combining an exponential term, representing the noise, and a Gaussian term, representing the signal peak, was employed to fit the spectrum in 5.4. The fitting result allowed for the extraction of the ADC resolution by calculating the ratio of mean value to the standard deviation from the Gaussian component: $\sigma/\mu = 76.9/189.7 = 40.6\%$. However, the Gaussian component was not able to perfectly reproduce the signal

peak because the peak signal did not come from a mono-energetic X-ray.

It is possible to obtain a pronounced escape peak in the ADC spectrum. To achieve this, the signals should be acquired from one multiplexing board, with the segment of 22×16 pads in proximity to the ^{55}Fe source, connected to the electronics while the rest of the pad plane and the micromesh are grounded. This configuration can increase the relative intensity of the escape peak to the K_α peak, thus making the escape peak more pronounced, as illustrated in Figure 5.5. A combinatory function of one exponential term and two Gaussian terms was used to fit the spectrum. The resolution of the signal peaks is 25.9% and 37.7% respectively.

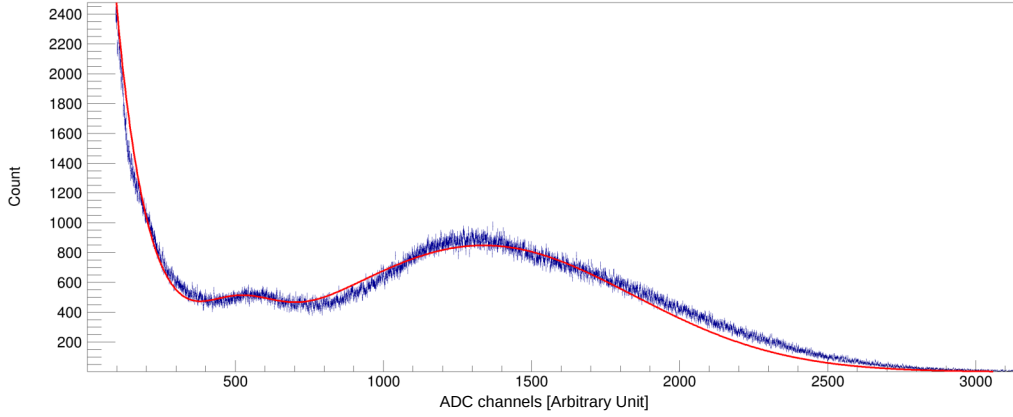


Figure 5.5.: X-ray ADC spectrum on a group of pads. The test duration is 302 s. The amplification gain is $20\times$. The blue line is the raw spectrum while the red curve is the fitting result.

Calculation of the primary current

Every signal count in Figure 5.4 represents one event of primary ionization whatever is the event. After subtracting the background noise using the exponential component from the fitting, the sum of all ADC bins (y_j) for the net signal peak can be calculated. Dividing the sum by the duration time t , the event rate can be subsequently obtained. Assuming, for simplicity, that all escape events result in full energy deposition and all incident X-ray energy is absorbed by the gas atoms and molecules, the primary current can be expressed as

$$I_{\text{primary}} = \frac{E_{\text{KX}}}{\epsilon_i} \cdot e \cdot \frac{\sum y_j}{t}, \quad (5.2)$$

where E_{KX} denotes the average energy of X-rays including K_α and K_β X-rays, e the electron charge. The W-value of the test gas mixture is $\epsilon_i = 28.8$ eV [232].

5.2. Gain Studies

5.2.1. Gain Measurements

The total effective gain of the entire amplification stage was determined. Utilizing the experimental setup described in Section 5.1.1 and the ^{55}Fe source, the current on the DLC, both with and without the source (I_{DLC} and $I_{\text{DLC_leakage}}$, respectively), was measured. The total effective gain of the detector can thus be calculated by

$$G_{\text{tot_eff}} = \frac{I_{\text{amplification}}}{I_{\text{primary}}} = \frac{I_{\text{DLC}} - I_{\text{DLC_leakage}}}{I_{\text{primary}}}, \quad (5.3)$$

where the total primary current, I_{primary} , is calculated from Equation 5.2 with summation over all signals on the micromesh. In the following sub-sections, the quantitative relationship between the total effective gain and either the electric field strength or the applied potential bias in each TPC region (Micromegas, transfer field, GEM, and drift field) are measured and analyzed. For each region, the electric field strength or potential bias was varied within a range covering the need for both the laser measurements and the final beam experiment, while maintaining constant values for the other regions. By extracting the total effective gain with Equation 5.3, its dependence on the potential difference or the electric field intensity in the various TPC regions was investigated.

Gain in the Micromegas

To assess the gain performance of the Micromegas, the second stage of the amplification region, measurements were conducted with the following HV settings:

$$\begin{aligned} E_d &= 220 \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 350 \text{ V}, \\ E_t &= 1000 \text{ V/cm}, \\ U_{\text{DLC}} &= [380, 440] \text{ V}, \end{aligned}$$

where the bias on the DLC layer, U_{DLC} , is incremented every 10 V.

As detailed in Section 2.2.3, the number of avalanche electrons generated per primary electron within the amplification region increases exponentially with the applied electric field. This relationship is confirmed by the total effective gain $G_{\text{tot_eff}}$ as a function of U_{DLC} which is the voltage difference between the DLC layer and the micromesh, as depicted in Figure 5.6. An exponential function is fitted to the data. In the logarithmic scale, the linear fit clearly demonstrates the exponential dependence of the gain on the voltage.

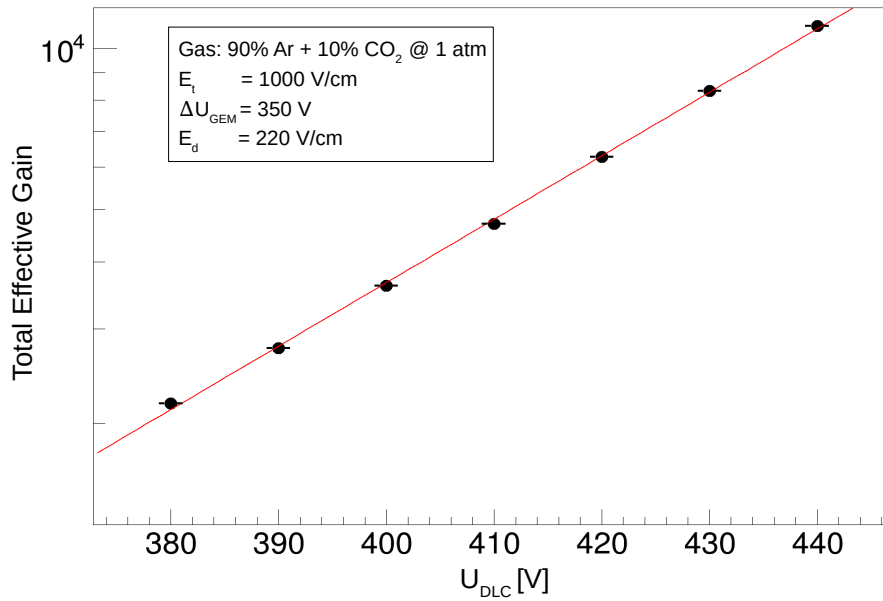


Figure 5.6.: Relation between the total effective gain $G_{\text{tot_eff}}$ and the high voltage bias on the DLC layer U_{DLC} . The vertical error bars are smaller than the marker size and therefore not explicitly shown. The red line is the linear fit of the data points in the logarithmic scale.

Gain as a Function of the Transfer Field

Within the transfer region, no amplification of the signal occurs. However, a stronger electric field E_t in this region allows for the more efficient extraction of electrons from the GEM into the Micromegas. Figure 5.7 showcases a result of gain measurements scanning over the transfer field with the following HV settings:

$$\begin{aligned} E_d &= 220 \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 310, 330, 350 \text{ V}, \\ E_t &= [200, 2400] \text{ V/cm}, \\ U_{\text{DLC}} &= 440 \text{ V}, \end{aligned}$$

where E_t is incremented every 200 V/cm.

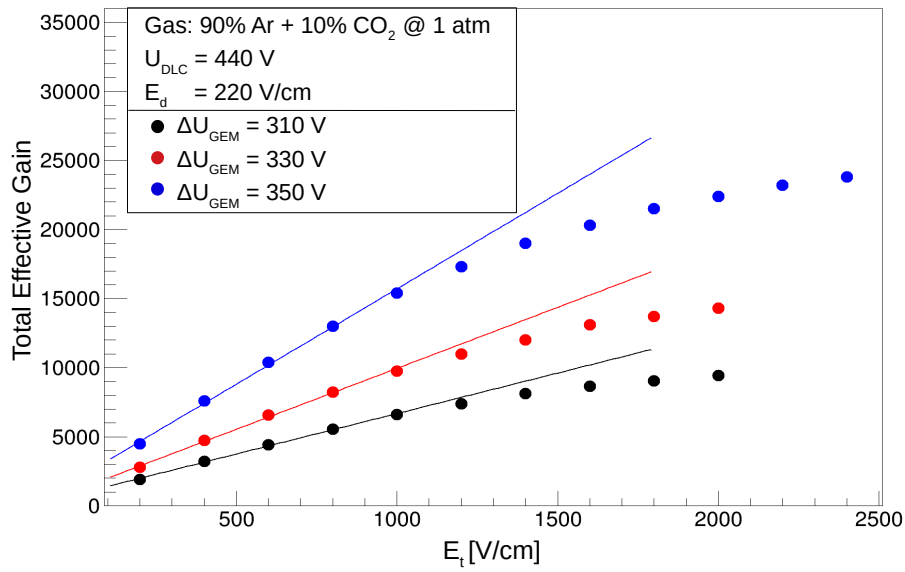


Figure 5.7.: Relation between the total effective gain $G_{\text{tot_eff}}$ and the transfer field E_t . The error bars are smaller than the marker size and therefore not explicitly shown. Lines of corresponding colors represent linear fits to the respective data points.

The gain $G_{\text{tot_eff}}$ as a function of the transfer field strength, presents a linear increase until approximately 1000 V/cm is reached. Beyond this value, the gain starts to saturate, indicating that the number of transferred electrons has reached its maximum limit. Notably, this behavior is independent of the voltage applied to the GEM. As shown in Figure 5.7, for all U_{GEM} values, the gain exhibits a linear increase from 200 V/cm to 1000 V/cm for the transfer field strength. Beyond 1000 V/cm, the saturation trend is consistent across all U_{GEM} values. This implies that the efficiency of electron extraction is not influenced by the quantity of electrons produced within the GEM stage. The phenomenon can be explained qualitatively as the following.

When the transfer field is zero, most of the electrons are absorbed by the bottom side of the GEM following the field lines instead of being extracted into the transfer field. In this case, the Micromegas mesh shares the same potential with the bottom side of the GEM but is more distant from the top side of the GEM, which leads to fewer field lines ending on the mesh or through it. With the increase of E_t , more field lines pass through the transfer region resulting in an increase of the gain. However, the extraction efficiency cannot be 100% due to the exact geometry. Some field lines inevitably terminate on the bottom side of the GEM leading to a loss of electrons, which depends solely on the high voltages applied onto the electrodes of the GEM as the Micromegas mesh is grounded.

Gain in the GEM

Similar to the Micromegas stage, the total effective gain increases exponentially with the voltage difference between the two sides of the GEM. Figure 5.8 illustrates for different transfer fields, the gain G_{tot_eff} as a function of the voltage on the GEM with the following HV settings:

$$\begin{aligned} E_d &= 220 \text{ V/cm}, \\ \Delta U_{GEM} &= [230, 430] \text{ V}, \\ E_t &= 0, 500, 1000, 1500, 2000 \text{ V/cm}, \\ U_{DLC} &= 440 \text{ V}, \end{aligned}$$

where the U_{GEM} is incremented every 10 V.

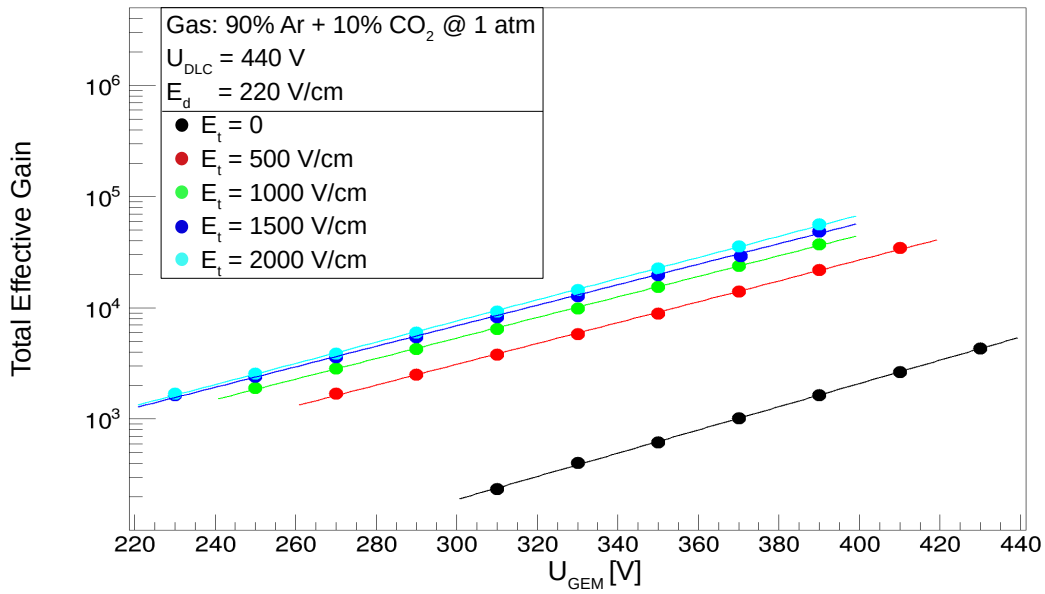


Figure 5.8.: The relation between the total effective gain G_{tot_eff} and the voltage applied on the GEM U_{GEM} . The error bars are smaller than the marker size and therefore not explicitly shown. Lines of corresponding colors represent linear fits to the respective data points in the logarithmic scale.

Analogous to the Micromegas, the total effective gain G_{tot_eff} displays a linear relationship with U_{GEM} in the logarithmic scale, as evidenced by the linear fits for each fixed E_t .

Notably, saturation of G_{tot_eff} is also observed in these measurements. Figure 5.8 indicates a saturation point for $E_t = 1000 \text{ V/cm}$. Beyond this value, G_{tot_eff} does not increase significantly with E_t . Nonetheless, this saturation does not affect the exponential dependence of G_{tot_eff} on U_{GEM} .

Gain as a Function of the Drift Field

While the drift field does not directly influence the amplification process, it can indirectly impact the total effective gain, G_{tot_eff} by affecting electron collection efficiency [233]. In weak electric fields, the probability of electron attachment is significantly increased. At extremely low field strengths, electron-ion recombination occurs. Furthermore, the increased diffusion of electron clusters in the weak fields can lead to electron absorption by GEM, resulting in electron losses and hindering their extraction from the drift region to the GEM's holes. Measurements of the relation between G_{tot_eff} and the drift field E_d were conducted with

following HV settings:

$$\begin{aligned} E_d &= [40, 220] \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 300, 350 \text{ V}, \\ E_t &= 1000 \text{ V/cm}, \\ U_{\text{DLC}} &= 440 \text{ V}, \end{aligned}$$

where E_d is incremented every 30 V/cm. Figure 5.9 illustrates a fraction of electron loss, for drift field strengths (E_d) below 100 V/cm. In this region, $G_{\text{tot_eff}}$ increases with E_d . Starting from $E_d = 100$ V/cm, the gain saturates, indicating that nearly all drift electrons are extracted by the GEM. It is noted that electron collection efficiency can be reduced from unity at extremely high drift field strengths (several kV/cm) [233]. However, due to the experimental setup's limitation of a maximum cathode voltage of 8 kV, the study of this effect under extremely high drift fields was not feasible.

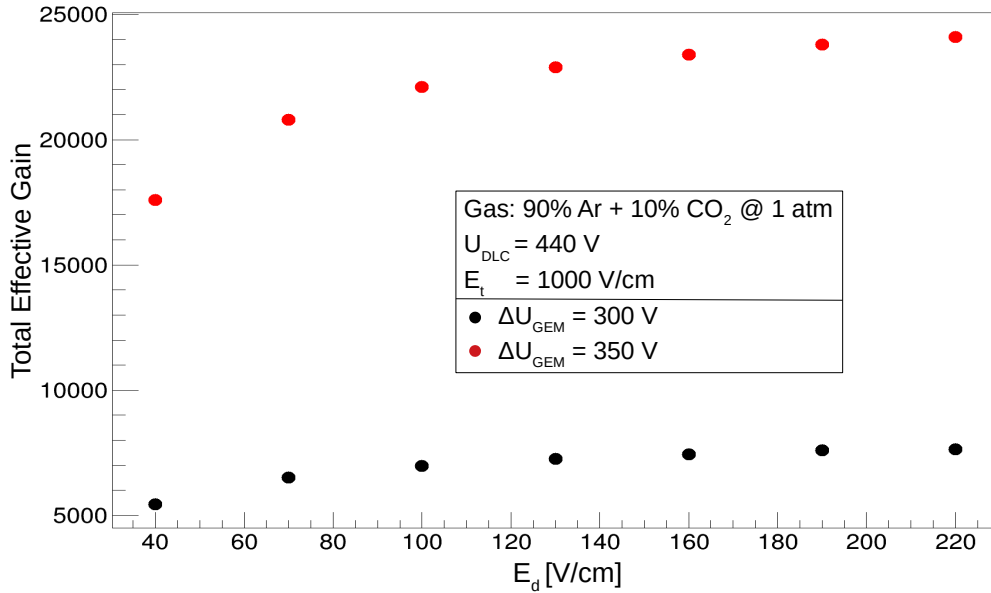


Figure 5.9.: The relation between the total effective gain $G_{\text{tot_eff}}$ and the drift field E_d . The error bars are smaller than the marker size and therefore not explicitly shown.

5.2.2. Discussion

The exponential dependence of the total effective gain, $G_{\text{tot_eff}}$ on the high voltage bias was assessed for the GEM and Micromegas stages of the HYDRA-TPC. Furthermore, the crucial role of efficient electron extraction from the drift field and transfer field to subsequent stages was demonstrated and quantified. These results are consistent with the expectations outlined in Section 2.2.3 and aligned with previous results in [234] with the same hybrid amplification configuration and the same gas mixture of 90% Ar + 10% CO₂.

Due to limitations in the experimental setup, specifically the inability to measure the cathode current, direct measurement of the ion back-flow (IBF) within the TPC has not yet been feasible. Consequently, the space charge effect in the drift region cannot be accurately assessed. This thesis adopts the IBF ratio of less than 1% for a gain of 5350, reported in [234] with a hybrid configuration of a GEM and a bulk Micromegas and the gas mixture of 90% Ar + 10% CO₂, as a reference value for the HYDRA-TPC. It was obtained by applying high

voltages with $\Delta U_{\text{GEM}} < 370$ V and the bias on the Micromegas, U_{MMG} , greater than 407.5 V. The resulting mean space charge density within the drift volume can reach a value of 50 fC/cm³ [210].

However, it is crucial to note that the IBF ratio and the resulting space charge determined in prior work do not accurately represent the conditions encountered in the in-beam experiments for the HYDRA-TPC. An accurate assessment of the IBF, depending on the effective total gain, and the space charge effect within the HYDRA-TPC requires systematic investigations under high-ionization conditions, such as measurements with high-intensity X-ray sources or heavy-ion beams. Furthermore, the measurement of the cathode current provides only an overall estimate of the space charge. To evaluate the distribution of the space charge and its influence on the tracking performance of the TPC accurately, in-beam assessments with reference laser tracks are necessary.

6. Laser Measurements with HYDRA-TPC

Measurements were performed using the HYDRA-TPC in conjunction with the laser system described in Section 3.3.1. These measurements aimed to evaluate the three-dimensional tracking capabilities of the TPC and investigate the impact of non-homogeneous magnetic fields within the GLAD magnet on electron drift. The general measurement procedure is outlined as follows:

1. Beam adjustment. By adjusting the beam angle from the source, the UV laser is directed into the active volume and subsequently ionizes the gas through the double ionization process.
2. Tracking. Hit information per event in the pad plane is extracted by analyzing the waveforms from the GET-based system. A tracking reconstruction method, incorporating the multiplexing scheme, is then employed to reconstruct the trajectories of the laser tracks within the TPC volume.
3. Inclusion of magnetic fields. To investigate the influence of magnetic fields on electron drift, non-uniform and non-parallel magnetic fields, created by the GLAD magnet, are introduced. The displacement of reconstructed tracks is analyzed under different magnetic field strengths.

The HYDRA-TPC was first tested with the laser system in the laboratory and subsequently operated within the GLAD magnet under various magnetic field strengths. Further details and results are presented in the following sections.

6.1. Laboratory Measurements

Preliminary test measurements of the entire detector system, integrating all components including the laser system and the GET-based system, were conducted in the local HYDRA laboratory. These initial tests aimed to verify the functionality of all elements and optimize the configuration of the GET-based system before laser measurements inside the GLAD magnet.

6.1.1. Test Setup

The test setup included the TPC, the high-voltage power supplies, the gas system, the GET-based system and the laser system.

High Voltage Modules and Settings

The high voltage configuration was identical to that described in Section 5.1.1, though with different power supplies. The DLC layer of the Micromegas was connected to a standalone Iseg power supply (the SHR series), while the mesh was grounded. This power supply supports a positive high voltage output of up to 6 kV. The remaining four electrodes, namely the cathode, the last strip, and both sides of the GEM, were connected to Iseg EHS 40 100n which supports negative outputs of up to 10 kV. The module was controlled by the master module Iseg CC24 and both were connected to an ECH-244 crate. In case of discharge, these four channels were configured to immediately shut down without a ramp-down process once a current exceeding the predefined limit was detected on any of the four channels during operation. The high voltage settings for

the test measurements were adjusted to achieve an overall amplification gain of approximately 1.4×10^4 , based on the gain measurement results presented in Section 5.2.1, as follows:

$$\begin{aligned} E_d &= 220 \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 400 \text{ V}, \\ E_t &= 1000 \text{ V/cm}, \\ U_{\text{DLC}} &= 400 \text{ V}. \end{aligned}$$

The leakage current stabilized at approximately 60 nA after the high voltage was applied to the Micromegas.

Gas System

During the laboratory tests, the alternative gas mixture of 92% Ar + 8% CO₂ served as the operating gas due to its availability. The flow of the gas mixture inside the TPC varied between 0.05 and 0.1 L/min during measurements, while 0.15 L/min during flushing. The pressure inside the TPC was not strictly controlled but kept slightly above the atmospheric pressure.

Electronics Configuration

The GET-based system combined with the multiplexing boards was connected to the TPC. There are mainly four parameters which were adjusted during the tests in the controller software of the GET-based system, namely the range, the peaking time, the writing clock frequency and the trigger delay:

1. The range parameter determines the maximum measurable signal range for the GET-based system, which corresponds to 4096 ADC channels. Four range options are available: 120 fC, 240 fC, 1 pC, and 10 pC. For all measurements presented in this thesis, the range was set to 120 fC to ensure precise readout of the waveform. Thus each ADC channel corresponds to a charge of approximately 183 electrons.
2. The peaking time determines the signal shaping time constant of the GET-based system. Possible values are: 171 ns, 232 ns, 280 ns, 568 ns, and several longer times up to 1 μ s.
3. The writing clock frequency determines the size of each unit time bin within the waveform, thus controlling the whole time window it covers. The duration of each time bin is fixed at $2^n \times 10$ ns, corresponding to writing clock frequencies of 100, 50, 25, and 12.5 MHz for $n = 0, 1, 2$, and 3, respectively.
4. The trigger delay defines the latency time for the GET-based system to initiate signal processing after receiving the trigger. The delay value is configured in the unit of 10 ns.

For both gas mixtures (90% Ar + 10% CO₂ and 92% Ar + 8% CO₂), the estimated minimum drift time for electrons to traverse the full cathode-to-GEM distance, based on drift velocities calculated using Magboltz, is around 20 μ s. Consequently, a writing clock frequency of 12.5 MHz was chosen to ensure complete coverage of the drift region and to allow for signal centering within the time window, which accordingly covers 40.96 μ s. However, for measurements with signals only from part of the drift region, a frequency of 25 MHz, corresponding to a time window of 20.48 μ s, is considered preferable to obtain a better time precision for the waveform.

The influence of trigger delay and peaking time on the waveform has been investigated through test measurements across different configurations. The results are discussed in the subsequent section.

6.1.2. Test Results

The laser source, upon activation, emits a beam directed into the active volume of the TPC, subsequently ionizing the gas in the active volume. Simultaneously the laser source transmitted a trigger signal to the

ZC706 boards at a frequency of 20 Hz for signal readout. After triggering, the GET-based system records the waveform within the predefined time window. To determine the best electronic configurations for waveform recording, several test measurements were performed with various GET settings.

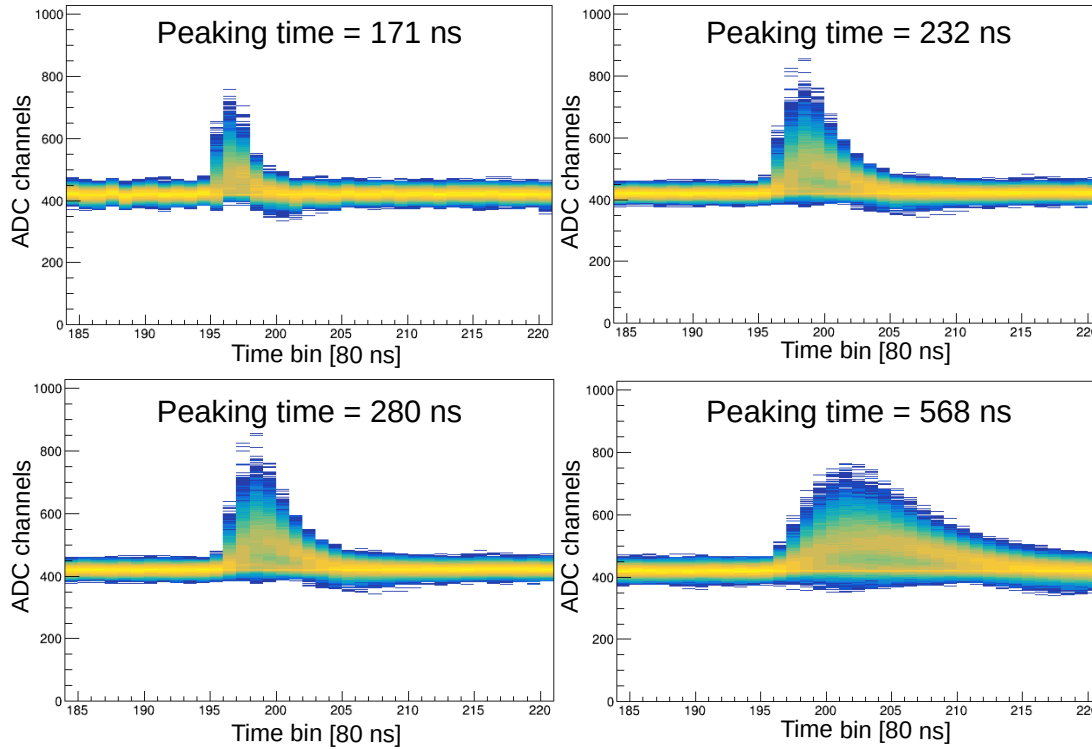


Figure 6.1.: Test measurement for different peaking times. For each peaking time, 500 laser events were taken from one AGET chip. Other common settings are: Writing clock frequency = 12.5 MHz; Trigger delay = 35 μ s; Range = 120 fC.

Peaking Time

The ion mobility for the 92% Ar and 8% CO₂ gas mixture is estimated to be approximately 2 cm/(V·s) based on Magboltz calculations. Consequently, a signal duration of approximately 200 ns can be obtained utilizing Equations 2.20 and 2.49. Given the available options in the GET-based system, a peaking time of 232 ns is considered optimal for integrating these signals.

Measurements with the laser were performed with various peaking times: 171 ns, 232 ns, 280 ns and 568 ns. The results are presented in Figure 6.1. A short peaking time of 171 ns in the GET-based system leads to loss of signal amplitude due to not enough time to fully integrate the signal as illustrated in the top left of Figure 6.1. On the other hand, a long peaking time of 568 ns results in a delay of the peak since the response is too slow compared to the rise of the signal as showcased in the bottom right of Figure 6.1. No significant performance difference was observed between peaking times of 232 ns and 280 ns. Given that 232 ns is closer to the 200 ns peaking time, it was selected as the optimal setting and consistently employed throughout the measurements.

Delay Time

Within a 40 μ s time window, signals from the entire active volume can be captured. Up to five distinct signal peaks can be observed in laser measurements, originating from the following sources:

1. Direct ionization of the gas by laser tracks from the micro-mirror bundles (as illustrated in Figure 6.2), corresponding to Peaks 1, 2, and 3.
2. Photoelectric effect on the GEM material due to laser reflection onto the bottom layer, corresponding to Peak 0.
3. Photoelectric effect on the cathode material due to laser reflection onto the cathode, corresponding to Peak 4.

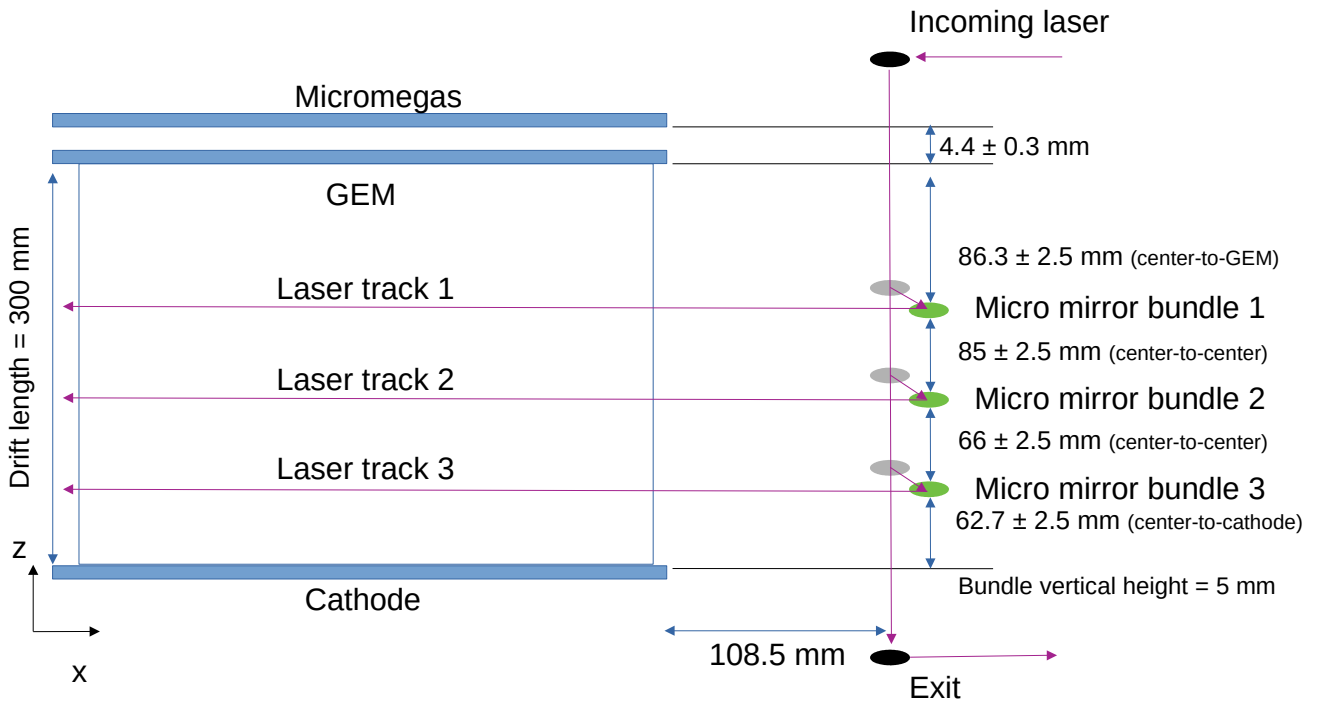


Figure 6.2.: Sketch of the sectional view of the TPC and the laser system (not to scale). The black ellipses represent reflection mirrors. The gray ellipses represent the beam splitters. And the green ellipses are the bundles of the micro-mirrors. "Center" in the figure refers to the central position of the micro-mirrors. The values are taken from the CAD drawings for the HYDRA-TPC.

Identifying all five peaks within the waveforms acquired by the GET-based system proved to be challenging, primarily due to the difficulty of precise laser beam alignment. The laser beam originates from a source positioned outside the TPC chamber and is projected into the active volume after undergoing multiple reflections including reflection by the micro-mirror bundles. Even a minimal deviation in the output angle of the laser from the source can significantly alter the individual peak amplitudes of the signals. This effect is particularly pronounced when the source is situated at a considerable distance from the TPC. Currently, the system lacks a precise mechanism for adjusting the output angle, relying solely on a manually operated knob on the source.

In addition, the chosen time window for data recording impacts the visibility and coverage of the laser peaks within the waveform. A systematic scan of the trigger delay was conducted across a range of 0 to 70 μs , given the writing clock frequency of 12.5 MHz. The trigger delay of 35 μs was found to effectively center all five peaks within the output waveform as illustrated in Figure 6.3 and was therefore selected in the laser measurements. The time of each peak in the waveform corresponds to the latency between the signal and the

trigger. Therefore, peaks at different time bins originate from various positions within the drift region. The five peaks in Figure 6.3 correspond to Peak 0 to 4, from left to right.

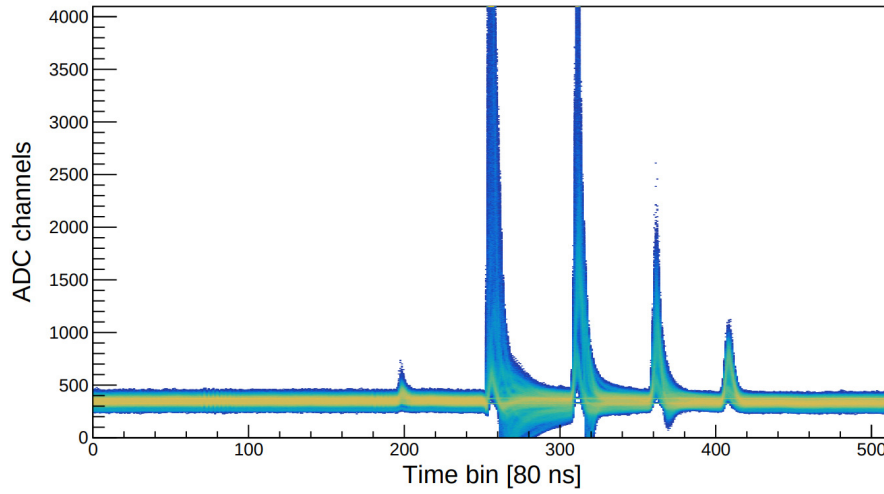


Figure 6.3.: Raw waveform for 1000 events from 64 channels of one AGET chip. From the left to the right, the five signal peaks correspond to Peak 0 to 4. The electronics settings are: Peaking time = 232 ns; Writing clock frequency = 12.5 MHz; Trigger delay = 35 μ s; Range = 120 fC.

Drift Velocity

This dependence of drift velocity and the drift field for the gas mixture of 92% Ar + 8% CO₂ was investigated as presented in Figure 6.4.

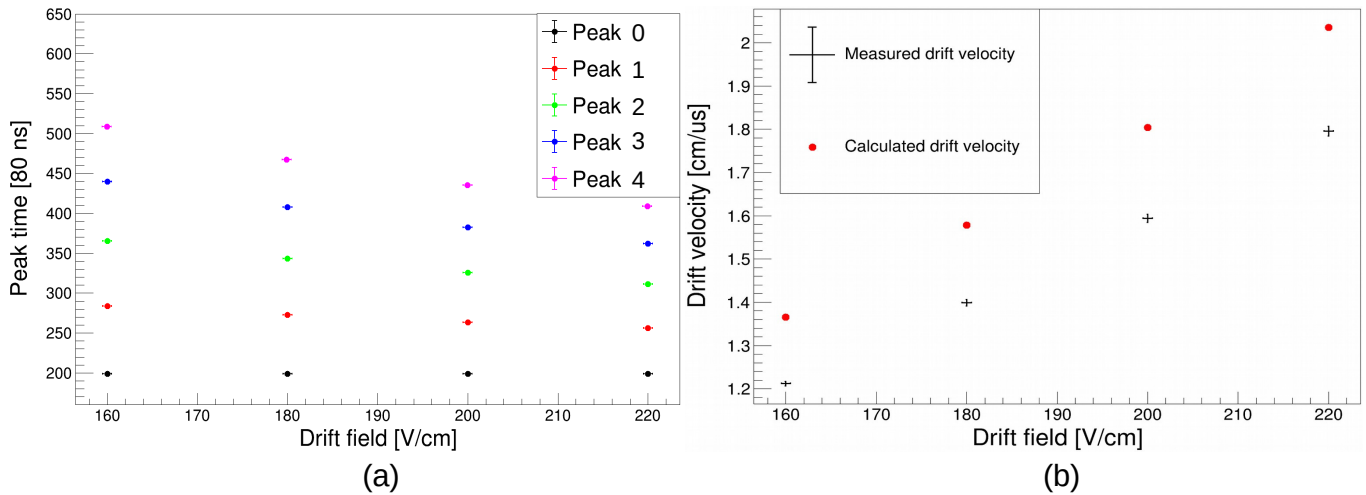


Figure 6.4.: (a) Peak time and (b) measured drift velocity as a function of the drift field. Calculated drift velocities are from [213], based on Magboltz. The electronics settings are identical to those for Figure 6.3.

The high voltage applied to the cathode and the last strip was manipulated to modify the drift field intensity. Subsequently, test measurements of the laser tracks were conducted under these varied drift field conditions.

To simplify the waveform analysis, the peak amplitude and its peak time are directly obtained from the raw waveform by identifying the maximum value and the corresponding time bin.

As illustrated in Figure 6.4(a), the position of the first peak in the waveform remains constant. This observation aligns with the expectation for signals originating from the GEM, where the arrival time of ionized electrons is independent of the drift field strength. Conversely, all other four peaks exhibit a downward shift in their peak positions with increasing drift field strength. Furthermore, the magnitude of this peak shift is proportional to the distance between the peak position and the GEM. The drift velocities for various electric fields were calculated based on the time difference between Peak 0 and Peak 4 corresponding to the entire drift distance of 30 cm as illustrated in Figure 6.4(b). The measured drift velocities exhibited a 0.1-0.2 cm/ μ s deviation above the theoretical values. This difference can be explained by a combination of factors, including a 5-degree higher gas temperature for the test measurement than the theoretical calculations and uncertainties in the gas composition of the test gas during production.

Peak Stability

To evaluate the stability of the laser peaks, a continuous one-hour laser measurement was conducted in the laboratory. This measurement aimed to observe the evolution of peak amplitude and time over time.

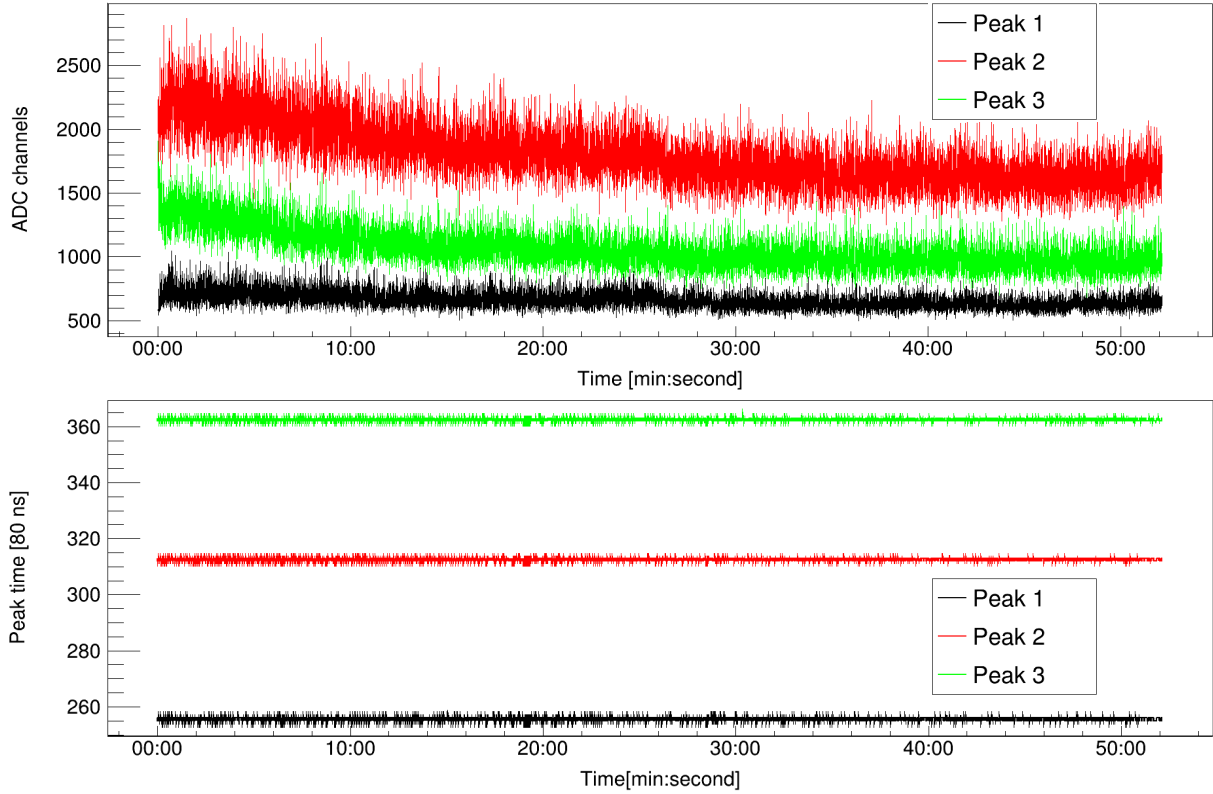


Figure 6.5.: Temporal evolution of (a) peak amplitudes and (b) peak time for three laser peaks. Peak amplitudes and time were continuously extracted from the waveform for Channel 26 of a tested AGET chip with the same electronics configurations as in Figure 6.3.

As depicted in Figure 6.5(a), the peak amplitude of the laser tracks exhibited a global decreasing trend, a phenomenon hereafter referred to as the quenching effect, with significant local fluctuations. For three specific peaks, the initial amplitudes in ADC channels were 690 ± 86 , 1919 ± 161 , and 1680 ± 221 , respectively. These values subsequently decreased to 662 ± 46 , 1531 ± 115 and 1062 ± 123 . The observed significant fluctuations

can be attributed to the low probability of the double-photon ionization process and gain fluctuations within the amplification region. For Peaks 2 and 3, the amplitudes decreased by approximately 400 and 600, respectively, over one hour. A plausible explanation for the quenching effect is instability in the beam angle emitting from the laser source and a misalignment of the incident laser beam. The underlying causes of this effect are currently under investigation.

The peak time, as depicted in Figure 6.5(b), remained stable, with mean values and standard deviations of $20.46 \pm 0.03 \mu\text{s}$, $24.92 \pm 0.04 \mu\text{s}$ and $28.93 \pm 0.04 \mu\text{s}$ for the respective peaks. This indicates a consistent gas drift velocity throughout the measurement.

6.2. In-GLAD Measurements

Two laser measurement campaigns with the HYDRA-TPC were conducted at Cave-C of GSI/FAIR, in November 2023 and February 2024. The November campaign, referred to as the first experiment in this thesis, aimed to commission and operate the entire setup within the GLAD magnet to assess the durability of all components under a high magnetic field of up to around 2 T. Additionally, the experiment investigated the influence of the magnetic field on laser tracks. The February campaign, referred to as the second experiment, aimed to provide more statistics for laser measurements and to preliminarily explore space charge effects within the TPC using the ^{12}C beam at an energy of 1.9 GeV per nucleon and the intensity varying from 10^3 to 10^6 particles per second. This thesis focuses on the first experiment. Sections 6.2.1 - 6.2.4 present the first experiment and the analysis results. Section 6.2.5 presents the second experiment and preliminary results.

6.2.1. Experimental Summary for the First Experiment

Detector Installation and Setup

The detector components were mostly the same as in the laboratory measurement in Section 6.1. The TPC was mounted inside the GLAD funnel. Its horizontality was verified using a spirit level. The laser source was positioned over 5 meters away from the TPC outside GLAD. The UV laser beam was emitted from the source at a frequency of 20 Hz, directed towards a reflection mirror, and subsequently reflected into the TPC as depicted in Figure 6.6.

The TPC was operated with a mixture of 90% Ar + 10% CO_2 under the pressure of around 10 mbar above the atmospheric pressure and a flow of 0.1-0.15 L/min. The general high voltage settings throughout the measurement are listed as follows:

$$\begin{aligned} E_{\text{drift}} &= 220 \text{ V/cm}, \\ \Delta U_{\text{GEM}} &= 380 \text{ V}, \\ E_{\text{transfer}} &= 1000 \text{ V/cm}, \\ U_{\text{DLC}} &= 400 \text{ V}. \end{aligned}$$

This configuration corresponds to approximately an effective gain of 10^4 based on the gain results in Section 5.2.1. The gain is lower than the value for the laboratory measurements in Section 6.1 to reduce the risk of signal overflow and potential discharge. The R³B cave was locked during operation to ensure the safety of personnel.

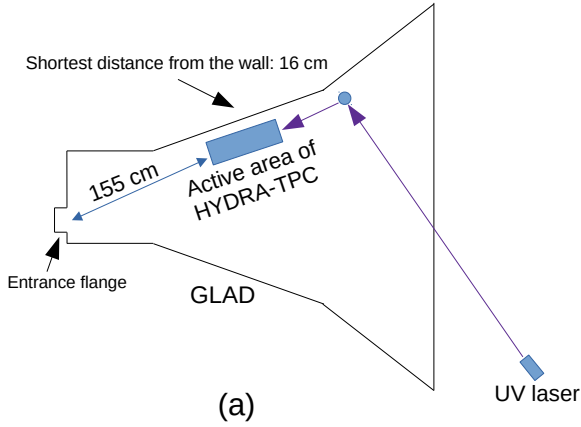


Figure 6.6.: (a) Sketch of the experimental setup for the first experiment. (b) Position of the HYDRA-TPC, the front-end electronics, and the reflection mirror in the GLAD magnet.

Electronics Configuration and Pad Plane Mapping

The front-end electronics, including 4 AsAd cards and 16 AGET chips of the GET-based system, were positioned downstream of the TPC in the GLAD magnet. The front-end electronics were connected to the TPC with shielded blue ribbon cables detailed in Section 3.3.2. In total, there were 1024 electronic channels which covered 5376 pads. Each AGET chip with 64 electronic channels was connected to a pad region of 22×32 pads via a multiplexing board. The mapping for connection is presented in Figure 6.7. Subsequently, the front-end electronics were connected to the ZC706 boards positioned outside GLAD at a distance of approximately 2 meters. During the measurements, the electronics and the electrical connections of cables stayed robust against the highest magnetic fields. The general settings of the GET-based system are summarized as follows:

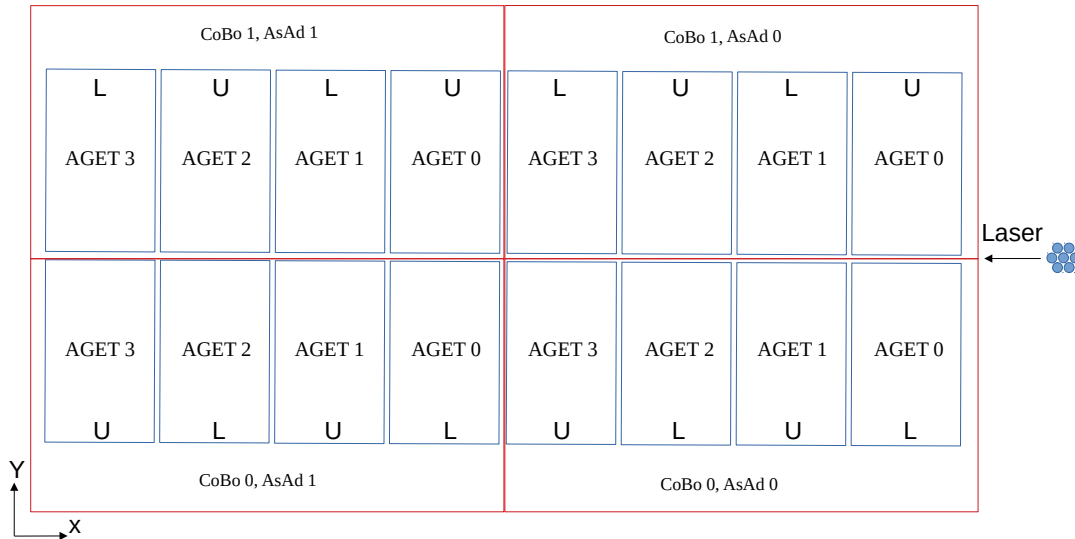


Figure 6.7.: Electronics mapping in the pad plane for the first experiment. 'U' and 'L' denote the sockets connection for the multiplexing boards in Figure 3.20.

Writing clock frequency = 25 MHz,
Trigger delay = 32 μ s,
Peaking time = 232 ns,
Range = 120 fC.

Data Collection

During the measurement, the current in the GLAD magnet was gradually increased from 0 to 3000 A, which corresponded to a magnetic field of 1.84 T, in several steps. A measurement run was conducted at each stage of the ramp-up process. Following several hours of measurement at the highest magnetic field of 1.84 T, the GLAD magnet was gradually ramped down to 0. Measurements were also taken during this ramp-down phase. The GET-based system was triggered by the laser source at a frequency of 20 Hz. Two CoBos were operated separately in a way that each time data were recorded by only one CoBo which covered half of the pad plane as showcased in Figure 6.7. The data collection for all measurement runs, with N_0 and N_1 representing the number of events recorded by CoBo 0 and 1, is detailed in Table 6.1.

Run Index	I [A]	B [T]	N_0	N_1	Others
1	0	0	5496	7596	
2	500	0.3	5642	5066	
3	1500	0.92	3935	21270	
4	1500	0.92	3935		$U_{DLC} = 420$ V, $\Delta U_{GEM} = 400$ V
5	3000	1.84		7900	$U_{DLC} = 420$ V, $\Delta U_{GEM} = 400$ V
6	3000	1.84	1470		$U_{DLC} = 415$ V, $\Delta U_{GEM} = 400$ V
7	3000	1.84	5474	1820	$U_{DLC} = 410$ V, $\Delta U_{GEM} = 400$ V
8	3000	1.84	3386	14942	
9	3000 to 1500	1.84 to 0.92		7930	
10	1500	0.92		3613	
11	1500 to 0	0.92 to 0		2406	
12	0	0	3915	14871	

Table 6.1.: Data for the first measurement campaign.

6.2.2. Waveform Analysis

Raw Waveform

Figure 6.8 presents superimposed waveforms from all channels within each AGET chip. Three distinct peaks were observed for AsAd 0 of CoBo 0, while no pronounced peak was found for AsAd 1. For CoBo 1, three distinct peaks were observed within the waveform time window for most AGET chips. However, only two peaks were observed for AsAd 0, AGET 0 and AsAd 1, AGET 3.

Before performing measurements, various electronic configurations were explored. This included employing a larger time window (Writing clock frequency = 12.5 MHz) and adjusting the trigger delay to values of 35 and 50 μ s. Despite these attempts, only the three aforementioned peaks were consistently observed at most. The issue was attributed to the increased difficulty of laser alignment due to the distance between the laser source and the TPC, as well as the reflection of the laser beam into the TPC.

Using the track reconstruction method outlined in Section 6.2.3, it was determined that the peaks from CoBo 1 do not form discernible tracks. Consequently, only tracks from CoBo 0 are analyzed in this thesis.

The three peaks from AGET 2, AsAd 0, CoBo 0 are located at 84.5 ± 1.8 , 241.2 ± 2.5 and 461.4 ± 3.2 in the

unit of time bin in Figure 6.8. Direct conversion of these positions to distances is not feasible since no absolute reference exists. The drift velocity of the gas mixture can deviate from the nominal value as already observed in the laboratory measurement (Figure 6.4(b)). Nevertheless, based on the ratio of relative distances between the peaks from the second experiment detailed in Section 6.2.5, it can be inferred that these three peaks correspond to Peaks 1, 2, and 4.

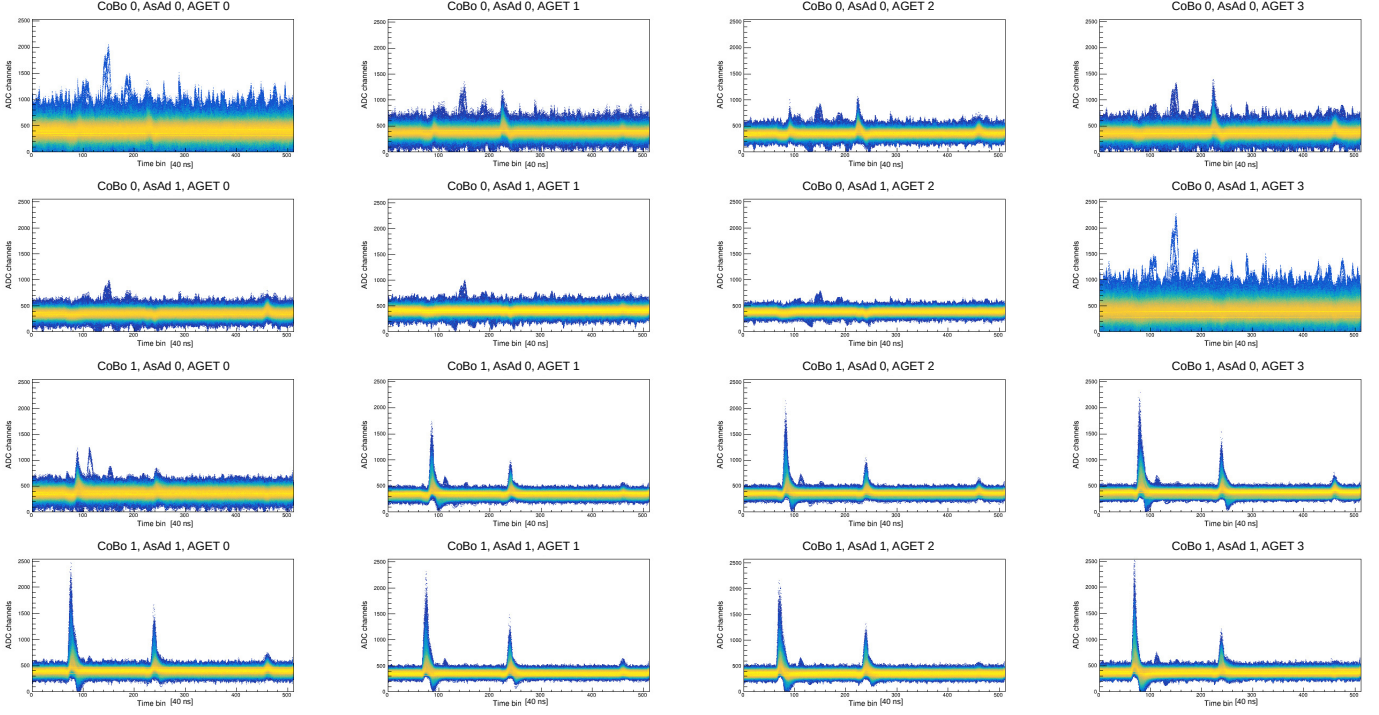


Figure 6.8.: Raw waveform for each AGET chip of both CoBo 0 and CoBo 1 from 1000 events of Run 1 in the first experiment.

Single peak amplitudes and time can be extracted from the waveform. The quenching effect was observed throughout runs of the first experiment while the peak time remained stable, as depicted in Figure 6.9. The peak amplitudes exhibited a faster quenching rate compared to the laboratory measurement (Figure 6.5). Peak 1 decreased from 761 ± 63 in ADC channels at 2:00 to 506 ± 19 at 5:00, Peak 2 decreased from 844 ± 117 at 2:00 to 521 ± 41 at around 8:00, and Peak 4 decreased from 568 ± 32 at 0:00 to 507 ± 27 at around 4:00. In Runs 4 to 8, no discernible signal peaks were observed, even with increased effective gain settings for the amplification region. Additionally, Peaks 1 and 2 in Figure 6.9 exhibited a unique behavior not observed in the laboratory measurement: a short period of amplitude increase from 731 ± 50 in ADC channels at 0:00 to 761 ± 63 at around 2:00 for Peak 1 and a rise from 736 ± 94 at 0:00 to a maximum of 844 ± 117 at around 2:00 for Peak 2, followed by the quenching effect. The underlying reason remains under investigation. The impact of the gas mixture or the electric field is considered insignificant as the drift velocity (represented by the positions of laser peaks) remained constant throughout the experiment, indicating no significant variation in gas properties or the uniformity of the drift field.

One potential explanation is a continuous shift with time in the incident laser during operation. This shift may originate from thermal effects arising from the laser's heating, leading to material expansion and subsequent deformation in components of the laser system, such as the support for the mirrors. Such deformation could

result in a deviation of the laser beam from its initial path, consequently changing the signal peak amplitude. Nevertheless, the orientation of the micro-mirrors remained stable, as evidenced by Figure 6.10. All three peaks on neighboring pads exhibit identical trends as in Figure 6.9, indicating that the direction of the sub-beams reflected by the micro-mirrors did not deviate from their initial paths within the active volume.

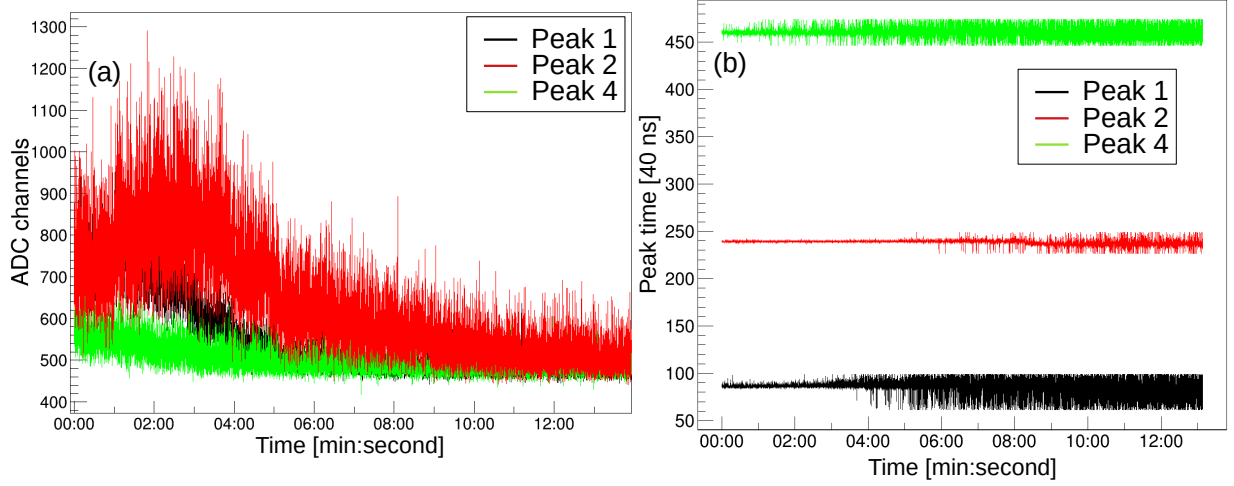


Figure 6.9.: Temporal evolution of (a) peak amplitude and (b) peak time for three laser peaks in Channel 26, AGET 2, AsAd 0, CoBo 1 from Run 3 of the first experiment.

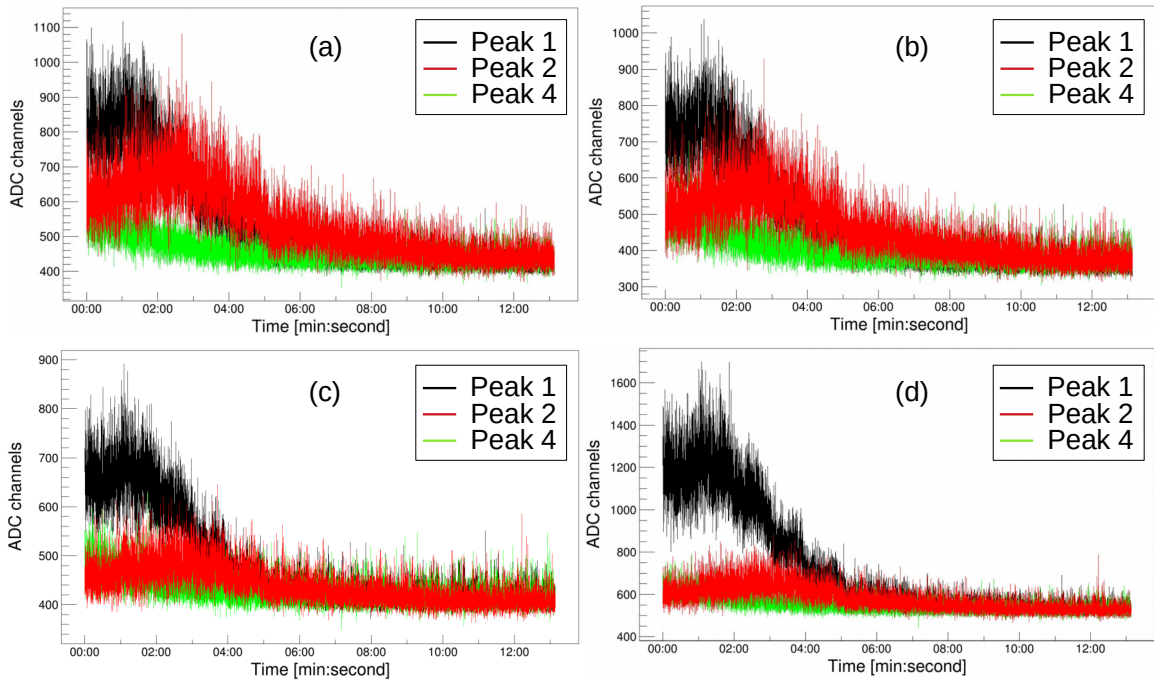


Figure 6.10.: Temporal evolution of peak amplitudes for (a) Channel 17, (b) Channel 21, (c) Channel 30, and (d) Channel 34 of AGET 2, AsAd 0, CoBo 1 from Run 3 of the first experiment. These channels correspond to the pads with a relative X index -2, -1, +1, and +2 in Figure 6.7, which are neighboring to the pad for Channel 26 according to the multiplexing scheme in Figure 3.20.

Noise subtraction and peak extraction

It can be noted that the significant fluctuations observed in Figure 6.8 (a) are due to the high noise level and the low probability of double ionization of the gas mixture by the UV laser. For further analysis, noise subtraction is necessary for the waveform.

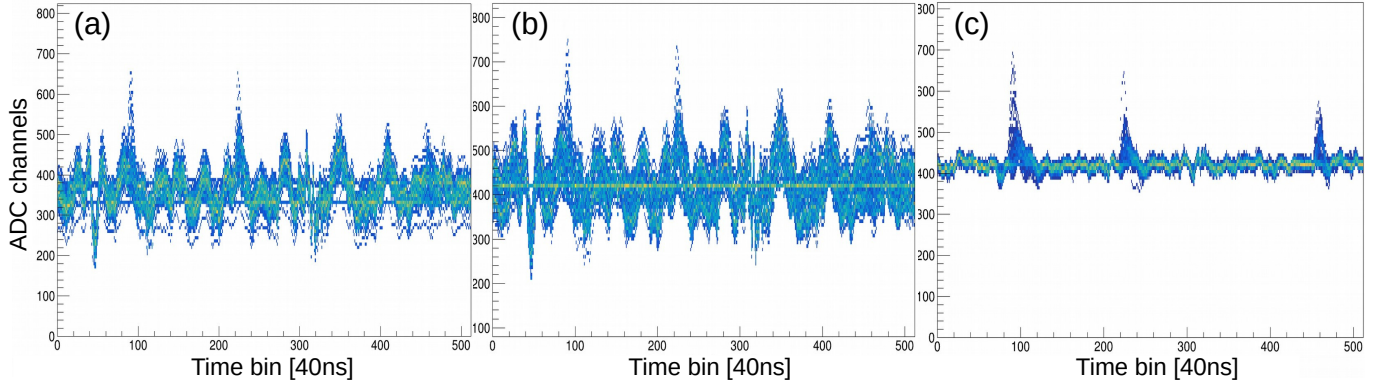


Figure 6.11.: Waveform of the 68 channels in AGET 2, AsAd 0, CoBo 0 of one event in Run 1 for the first experiment. (a) Raw waveform; (b) Waveform after FPN correction; (c) Waveform after noise subtraction.

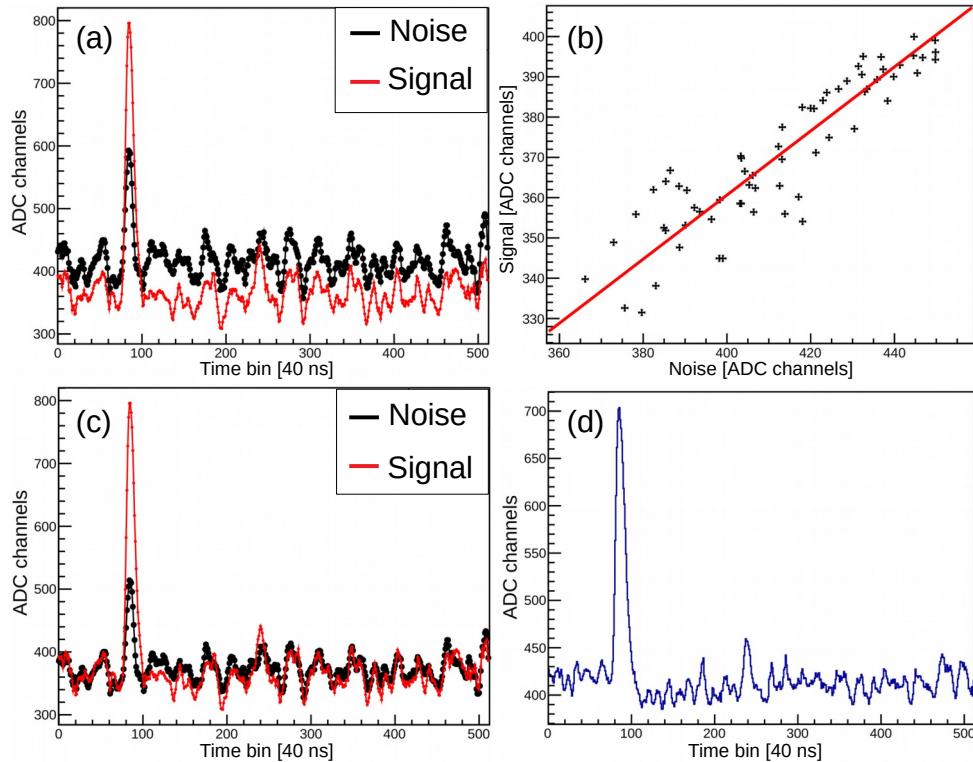


Figure 6.12.: Noise subtraction method. (a) Waveform of the noise and the signal channels; (b) Linear fitting of the relation between the noise and the signal in the time bin range [0,65]; (c) Waveform of the noise channel after scaling and shifting and the signal channel; (d) Waveform of the signal channel after correction.

An AGET chip includes 4 dedicated Fixed Pattern Noise (FPN) channels, each in every 17 channels. These FPN channels are not connected to any input signal, serving the specific purpose of measuring the electronic noise within the system. The waveform from each FPN channel (presented as approximately flat lines in Figure 6.11(a)) can be subtracted from the corresponding waveforms of other channels on a time bin-by-bin basis for each event. These FPN channels may exhibit baseline variations. Therefore, an equalization step is necessary to ensure consistency of baseline among channels. The FPN correction process for each waveform is applied as follows:

1. Identify the corresponding FPN channel;
2. Subtract ADC values for each time bin of the FPN channel waveform from the corresponding values in the specific channel waveform;
3. Apply an offset correction: calculate the difference between the mean value of the FPN channel waveform and 420 (the average ADC channels of the FPN channels over all AGET chips), and subsequently add this value to the waveform obtained from step 2, such that the baselines of all channels are aligned to 420.

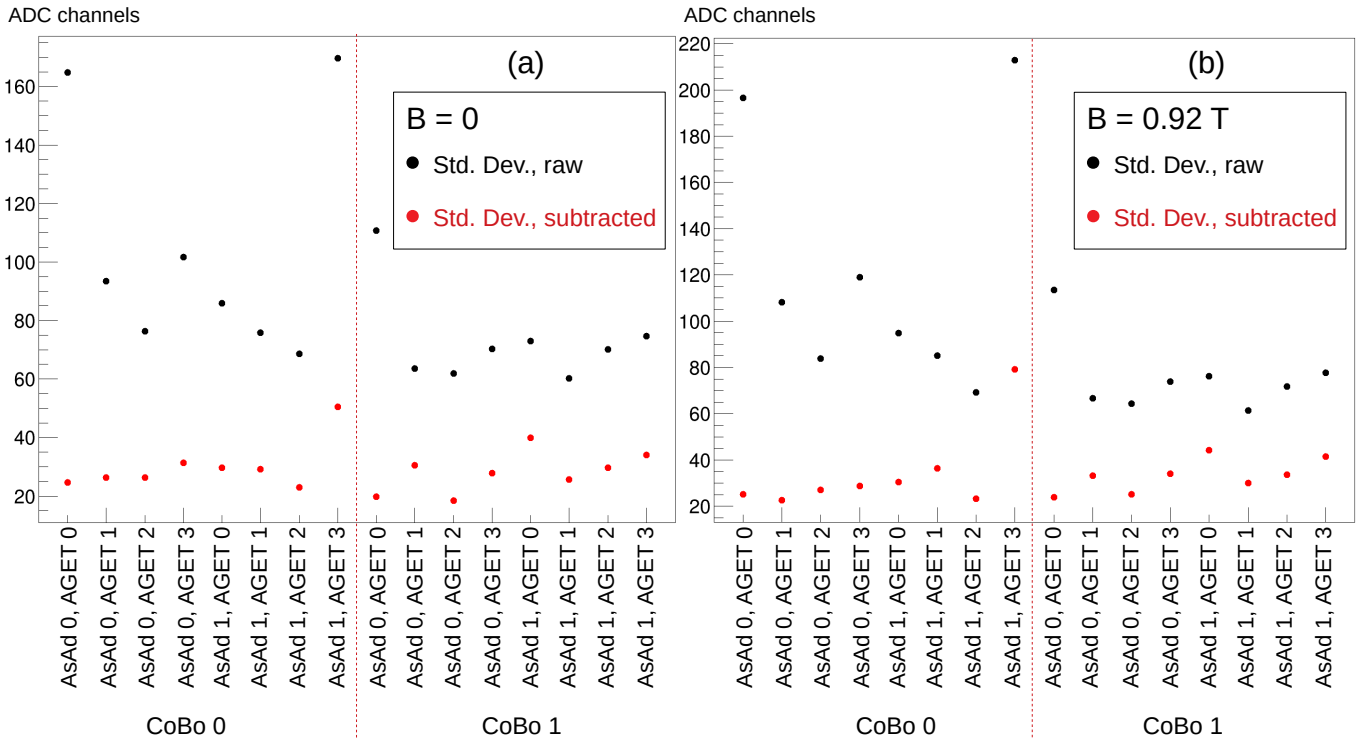


Figure 6.13.: Calculated standard deviation of the baseline in all AGET chips for (1) Run 1 and (2) Run 3. The black line represents the values before noise subtraction, while the red line represents the values after noise subtraction. The error bars are smaller than the marker sizes and therefore not explicitly shown.

The resulting waveform after FPN correction is presented in Figure 6.12(b). After FPN correction, the standard deviation in the baseline regions where no signal is expected (time bin range: [0,70], [110,210], [240,450] and [470,512]) of the waveform in Figure 6.11(b) is 56.9 ± 0.1 in ADC channels. It is noticeable in Figure 6.11(b) that within this AGET, all waveforms share a similar noise pattern. The origin of the noise has not yet

been identified. It may arise from the cross-talk between wires on the multiplexing boards or electromagnetic interference coupled into the electronics system due to inadequate shielding. A noise subtraction method is therefore employed as illustrated in Figure 6.12:

1. By scanning through the waveform within one AGET chip, select a ‘noise channel’ exhibiting the minimum peak amplitude at the signal position, since electromagnetic interference may induce peaks in noise channels;
2. Compare the waveform of a ‘signal channel’ and the ‘noise channel’ as plotted in Figure 6.12(a);
3. Within the baseline range (corresponding to time bin range: [0,65]), fit the relationship between the signal channels and the noise channels with a linear function as illustrated in Figure 6.12(b);
4. Extract the fitting parameters, the slope a and the intercept b , from the fitted linear function;
5. Subtract the noise from the raw signal (Figure 6.12(c)) with the fitting parameters:

$$ADC_{\text{signal,corrected}} = ADC_{\text{signal}} - (a \cdot ADC_{\text{noise}} + b). \quad (6.1)$$

With the above method, the noise level can be suppressed as illustrated in Figure 6.12(d). The standard deviation of all waveforms within a single AGET chip, measured in the baseline region of Figure 6.11(c), is reduced from 56.9 ± 0.1 ADC channels in Figure 6.11(b) to 12.54 ± 0.03 ADC channels after correction. This corresponds to an ENC of 2500 electrons, which is comparable to the ENC values with the cable in Table 3.1. The noise subtraction method can also be applied to other AGET chips and runs in the first experiment. For instance, Figure 6.13 summarizes the calculated standard deviation of the noise for Runs 1 and 3.

6.2.3. Laser Track Reconstruction

Following noise subtraction from the signal waveform, the maximum value above the baseline in each peak region can be extracted from the waveform as the peak amplitude for each channel of each AGET chip. The peak time corresponds to the time bin of the maximum value. Subsequently, a common threshold is employed to filter out peaks with amplitudes near the noise level and exceptionally high noise. This process generates cleaner events. The peak amplitude, represented by the difference between the absolute peak amplitude and the baseline in ADC, corresponds to the charge collected within the channel, such that a peak exceeding the threshold is regarded as a hit in the pad plane. Information of all hits in one event can be obtained for each pad based on the mapping between AGET chips and their corresponding locations in the pad plane from Figure 6.7 and the multiplexing scheme in Figure 3.20. Figure 6.14 presents all hits in the pad plane for Peak 1 on CoBo 0 from Run 1 with average amplitude for all hits. Subsequently, the laser tracks can be reconstructed from all these hits using the de-multiplexing method in the following. Using this track reconstruction method, it was concluded that the laser tracks for Peaks 1 and 2 on CoBo 0 can be reconstructed for the first experiment. In contrast, no discernible track could be reconstructed from the peaks on CoBo 1.

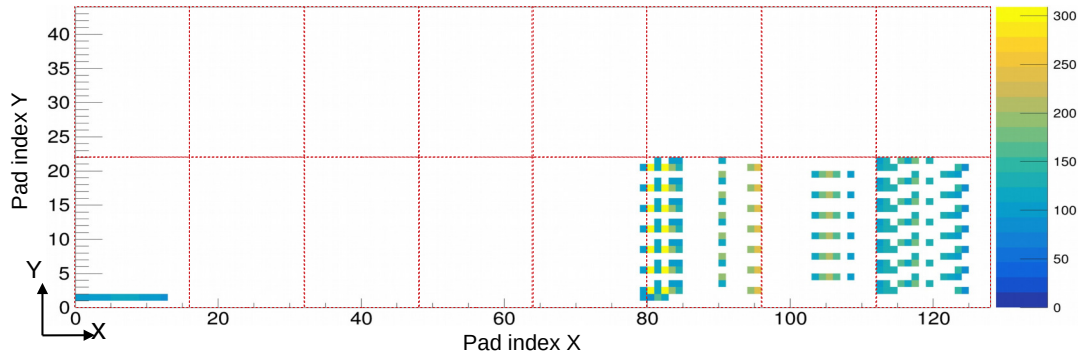


Figure 6.14.: All hits in the pad plane from one event for Peak 1 from Run 1 on CoBo 0 in the first experiment. The color denotes the signal amplitude in ADC channels above the baseline. The threshold is 90 in ADC channels above the baseline. Red dashed rectangles represent the area of each AGET chip detailed in Figure 6.7.

De-Multiplexing Method

According to the multiplexing scheme (Figure 3.20), every fourth channel of a multiplexing board is not multiplexed. These non-multiplexed channels are connected to the outermost row of pads on each side of the pad plane. These non-multiplexed pads thus can provide precise information about exit hits at the edge of the pad plane, serving as the basis for the de-multiplexing process.

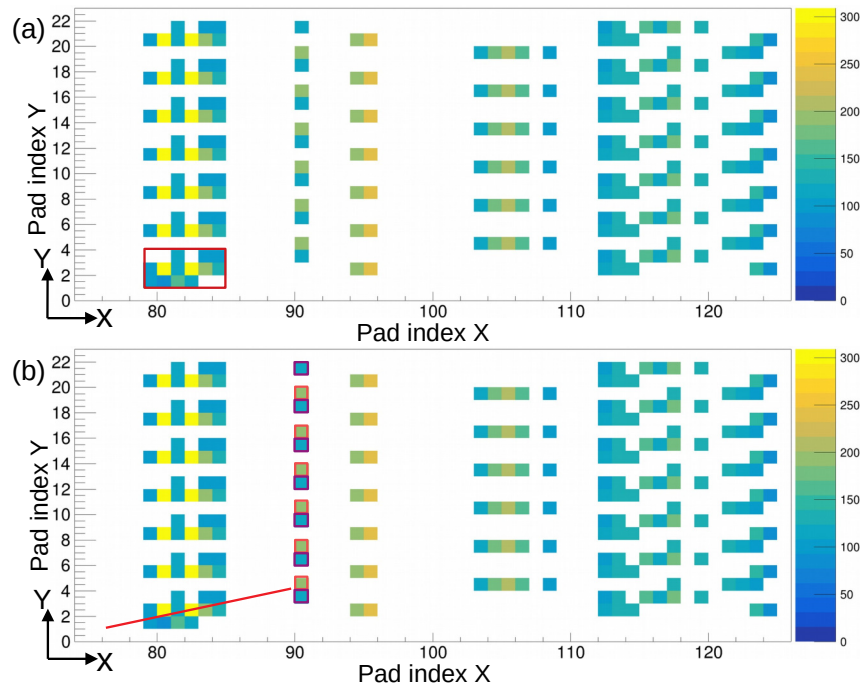


Figure 6.15.: De-multiplexing method: (a) Selection of non-multiplexing pads and their neighboring multiplexed pads in the red rectangle; (b) Selection of the minimum-distance pads from the multiplexing group. The color of the boxes denotes the peak amplitude above the baseline in ADC channels. The red line is from the previous fitting. Red or yellow squares are from the same multiplexing group.

By exploiting this information, a simple de-multiplexing method can be established as follows:

1. Search for all hits in non-multiplexed pads (The pad's Y index, $j = 1$ in Figure 6.14);
2. Along the X axis, calculate the centroid of charge $\bar{Y}_i$ for each X index in the non-multiplexed pad (as illustrated in Figure 6.15(a)) and the first group of multiplexed pads calculated by

$$\bar{Y}_i = \frac{\sum Q_{ij} \cdot Y_{ij}}{\sum Q_{ij}}, \quad (6.2)$$

where the index i and j denote the X and Y indices of the pad, respectively; Y_{ij} and Q_{ij} denote the Y coordinate of and the ADC channels of Pad (i,j) , respectively;

3. Perform a linear fitting to determine the relation between these centroids $\bar{Y}_i$ and their Pad Index X_i ;
4. Extrapolate the linear function from the fitting beyond the cluster with the non-multiplexed pads;
5. At Index X_i , select those pads belonging to the same multiplexing groups that exhibit the minimal distance to the line from the fitting in the previous steps, as illustrated in Figure 6.15(b);
6. Calculate the centroid of charge $\bar{Y}_i$ at Index X_i of these pads;
7. Perform another linear fitting including both the new centroid of charge $(X_i, \bar{Y}_i)$ and those previously identified;
8. Repeat Steps 5-7 until the linear function reaches the last X index with hits as illustrated in Figure 6.16(a).

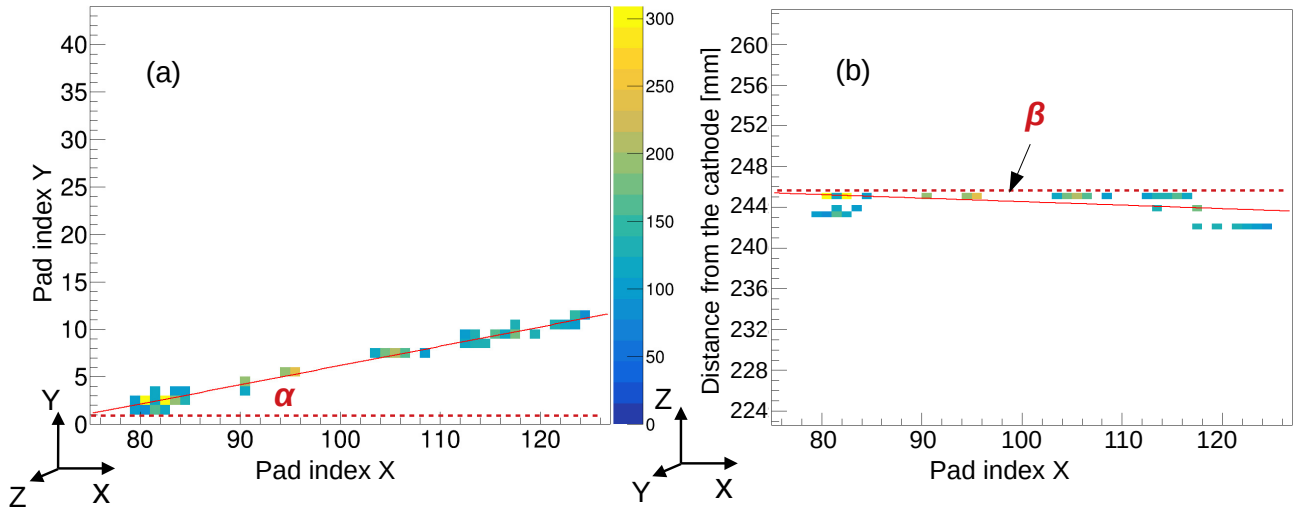


Figure 6.16.: Track reconstruction in (a) the XY plane and (b) the XZ plane from all hits in Figure 6.14. The color of the boxes denotes the peak amplitude above the baseline in ADC channels. α denotes the angle between the reconstructed track and the XZ drift plane. β denotes the angle between the reconstructed track and the XY pad plane. The red solid lines represent linear fits to the hits.

It is essential to note that the de-multiplexing method described above is strongly dependent on the presence of hits on the non-multiplexing pads. Without any hit on these pads, the track reconstruction procedure cannot

be initiated. Furthermore, events with a limited number of hits on these pads can lead to significant dispersion in the initial linear fit during Step 3. Consequently, the extrapolation of this linear function to identify multiplexed pads in Steps 4 and 5 may result in incorrect pad correlation and subsequently deviate from the true track. Figure 6.17 showcases the distribution of linear fit parameters for 437 events, all reconstructable tracks in Run 1. The observed inverse correlation between the intercept and slope of these fits demonstrates the method's reliance on non-multiplexed pads. Additionally, in Figure 6.17, there are a few scattered events with slopes less than 0.14, attributed to false correlations during Steps 3 to 5 of the track reconstruction process. To enhance the performance of the track reconstruction, an additional constraint, within the range of 0.18 to 0.22, is imposed a posteriori on the slope during every step of the fitting.

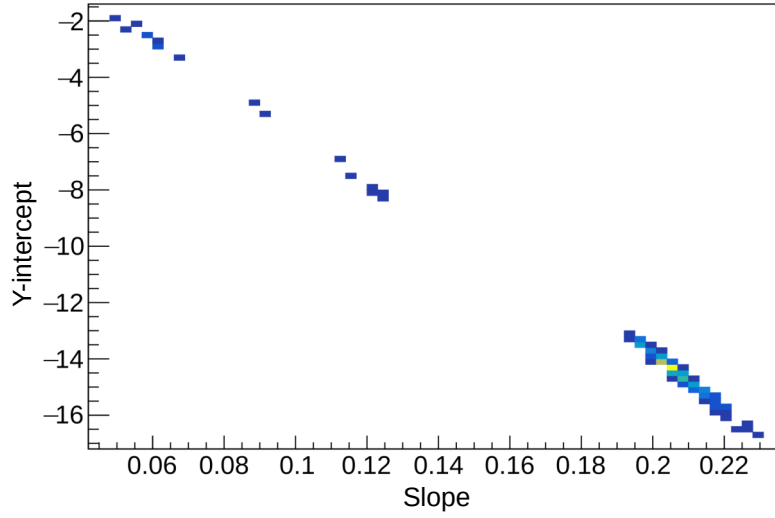


Figure 6.17.: Fitting parameters of the linear functions for all reconstructed tracks of Peak 1 from Run 1 on CoBo 0.

The reconstruction of the track above is performed in the pad plane which is merely a projection of the real laser track in the XY plane. With selected hits, the projection of the laser track in the XZ drift plane can be reconstructed as illustrated in Figure 6.16(b). The distance from the cathode for each hit is calculated by multiplying the peak time and the drift velocity calculated in the second experiment detailed in Section 6.2.5. The linear functions resulting from the track fittings in both the XY pad plane and XZ drift plane are plotted in Figure 6.16. The fitting parameters, the slopes b_{XY} and b_{XZ} of the linear functions for the tracks in the XY plane and the XZ plane respectively, can be subsequently interpreted to determine the angles of the laser track. The angles α and β defined in Figure 6.16 can be calculated by

$$\alpha = \arctan(b_{XY}), \quad (6.3)$$

$$\beta = \arctan(b_{XZ}). \quad (6.4)$$

The angular distributions of α and β for the reconstructed tracks of Peak 1 are presented in Figure 6.18(a). While the third peak (Peak 4), originating from the cathode, does not present any track, the same track reconstruction procedure can be applied to the second peak (Peak 2) with the angular distributions showcased in Figure 6.18(b). There are in total 2162 tracks reconstructed for the second peak in non-magnetic field measurements with a threshold of 140 in ADC channels. For Peak 1, the mean values for the angles α and β are calculated to be $11.74^\circ \pm 0.36^\circ$ and $0.00^\circ \pm 0.77^\circ$, respectively. For Peak 2, the mean values for the angles α and β are determined to be $9.69^\circ \pm 0.17^\circ$ and $-1.20^\circ \pm 0.35^\circ$, respectively. These angles from

reconstruction, however, do not correspond to any measured orientation angles of the mirco-mirrors before the bundles were mounted inside the TPC. One possible explanation is that due to the installation process, the angles of micro-mirrors were altered.

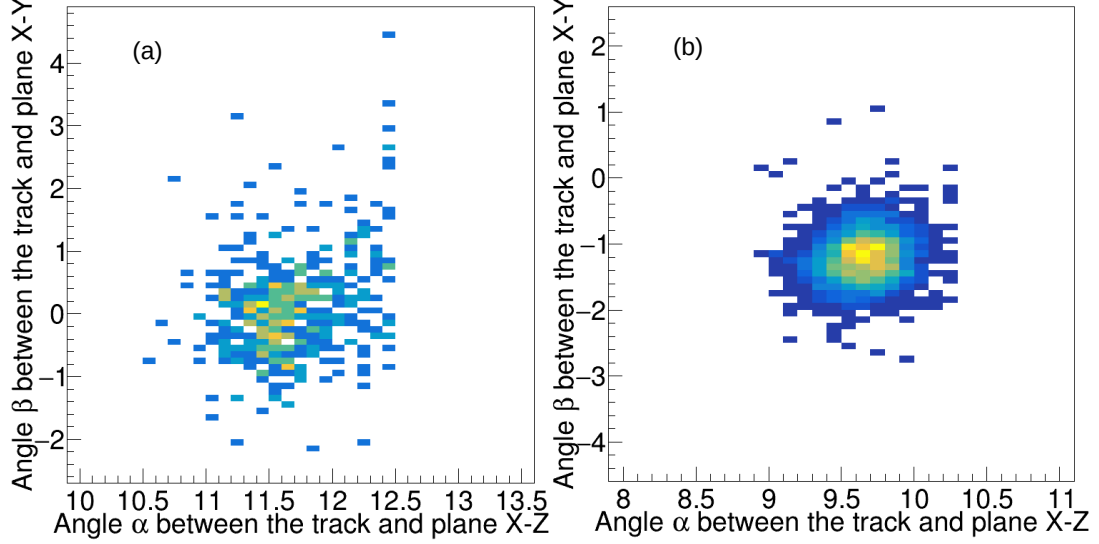


Figure 6.18.: Angular distribution of the reconstructed tracks for (a) Peak 1 and (b) Peak 2 on CoBo 0 in Run 1. The angles are plotted in degrees.

To assess the effectiveness of the track reconstruction method, the fitting functions are extrapolated to the Pad X index corresponding to the micro-mirror's position, as illustrated in Figures 6.2 and 6.7, to determine the associated Y coordinate of the micro-mirror. Figure 6.19 presents the distribution of the micro-mirror's Y coordinates for Peak 1 and 2 from extrapolation, with mean values and standard deviations of $46.3 \text{ mm} \pm 1.1 \text{ mm}$ and $42.5 \text{ mm} \pm 0.6 \text{ mm}$, respectively. Given the uncertainty regarding the specific micro-mirror within the bundle responsible for reflecting the laser track of Peaks 1 and 2, and considering that the micro-mirror bundle center is aligned with the pad plane center, the theoretical Y-coordinate of the micro-mirror is given by $44 \pm 1.5 \text{ mm}$. Both reconstructed laser tracks coincide with the position of their respective micro-mirrors within the estimated uncertainty.

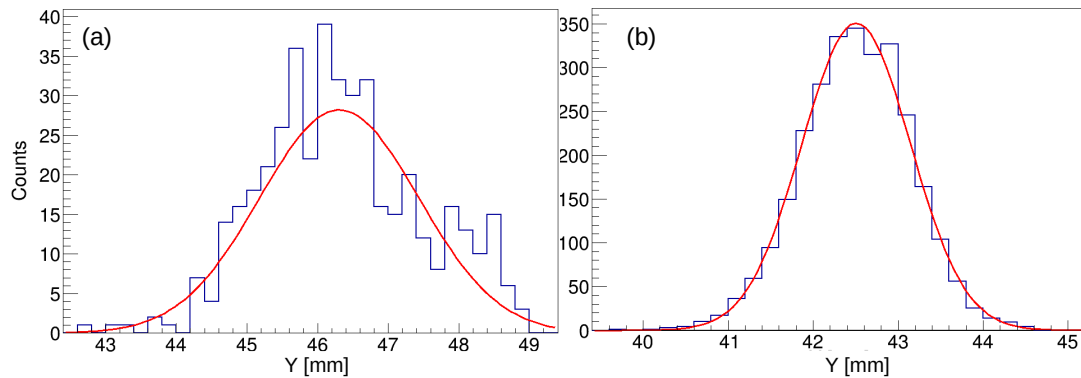


Figure 6.19.: Extrapolated Y positions of the micro-mirror from track reconstruction for (a) Peak 1 and (b) Peak 2 in Run 1 on CoBo 0.

The spatial resolution for the laser track can also be extracted by calculating the dispersion of the centroid of charge from the reconstructed track. The minimum distance D_{ij} between any centroid of charge (x_i, y_j) and the track $y = p_0 \cdot x + p_1$ is given by

$$D_{ij} = \frac{|y_j - p_0 \cdot x_i - p_1|}{\sqrt{p_0^2 + 1}}, \quad (6.5)$$

where parameters p_0 and P_1 are determined from the mean values of the fitting parameters obtained by averaging over all reconstructed laser tracks. The distances from all charge centroid positions to these lines were computed with Equation 6.5 and are presented in Figure 6.20.

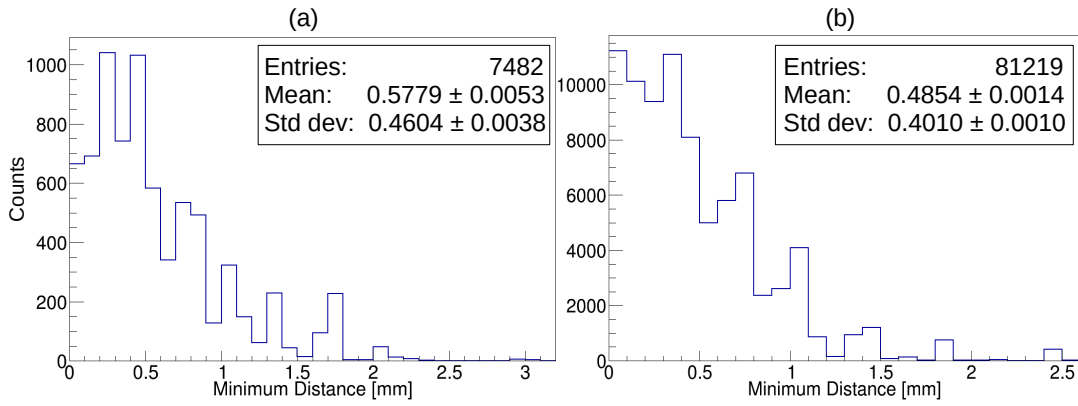


Figure 6.20.: Distances distribution of centroids of charge to reconstructed laser tracks in the pad plane for (a) laser Peak 1 and (b) Peak 2 from Run 1 in the first experiment.

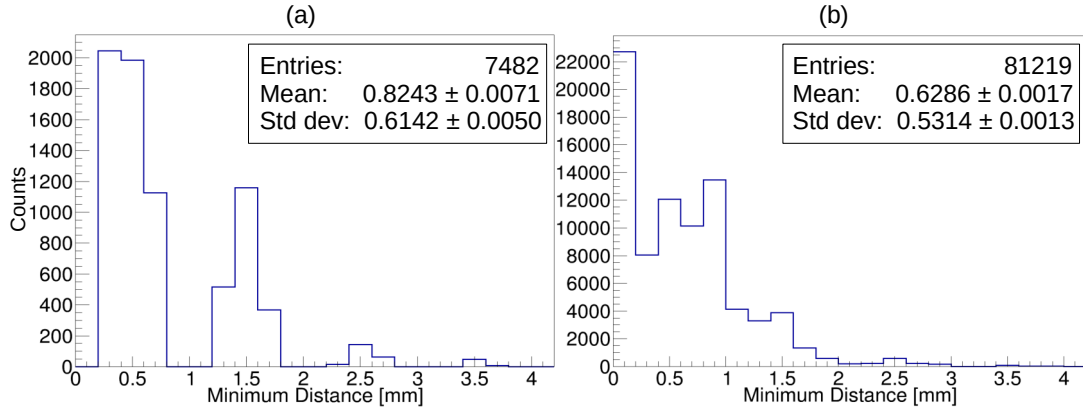


Figure 6.21.: Distances distribution of selected hits to reconstructed laser tracks in the drift plane for (a) laser Peak 1 and (b) Peak 2 from Run 1 in the first experiment.

The mean distances of 0.578 ± 0.005 mm and 0.485 ± 0.002 mm were obtained for Peak 1 and Peak 2, respectively. Given that the maximum laser beam radius, equivalent to the radius of one micro-mirror, is 0.5 mm, the criterion of $1\text{-}\sigma$ spatial resolution of $200\text{ }\mu\text{m}$ is expected to result in a dispersion of 0.54 mm. Under a $1\text{-}\sigma$ criterion, Peak 2 meets the expected dispersion, though Peak 1 slightly exceeds it.

The same method was employed to calculate the spatial resolution in the drift plane (XZ). Figure 6.21

illustrates the distribution of distances between selected hits associated with their corresponding peak times and the reconstructed track. The mean distances of 0.824 ± 0.007 mm and 0.629 ± 0.002 mm were obtained for Peak 1 and Peak 2, respectively. Based on a $1\text{-}\sigma$ criterion, both peaks demonstrate a spatial resolution better than the 3 mm requirement in the drift plane.

Nevertheless, the de-multiplexing process has inherent flaws, which resulted in degraded spatial resolution, as elaborated in the following:

1. A range calibration test for the GET-based system is not yet available. As illustrated in Figure 6.14, certain pads exhibit low peak amplitude values or are below the threshold along the laser track path. A potential explanation for this discontinuity of the track is that the range is not uniform across all channels for the electronics, despite identical range settings within the GET software controller. As a consequence, uncertainty arises in the calculation of the charge centroid and subsequent fitting process. To correct the amplitude difference, a calibration procedure for the GET-based system is desired.
2. A bias is introduced in the centroid of charge calculation when using non-multiplexed pads. Due to the proximity of the pads at the edge of the active area, a fraction of the charge may not be collected by the electronics. This missing charge information leads to a systematic shift of the calculated centroid towards the center of the pad plane in Step 2 of the method.
3. The centroid of charge calculation only on Y_i in Equation 6.2 introduces additional systematic uncertainty. Charge clusters exhibit diffusion in all spatial dimensions, which can extend beyond the area of a single pad and are particularly pronounced in the presence of the resistive layer of the Micromegas. As a result, a single charge cluster can induce signals on multiple pads. To simplify the de-multiplexing process, Step 2 of the method restricts charge diffusion to the Y-axis for centroid calculation, assuming confinement within a single pad along the X-axis. This simplification deviates from the true physical behavior. In addition, it degrades the spatial resolution along the X-axis, introducing an uncertainty equivalent to half the pad size.

In the next phase of the HYDRA program, the implementation of the non-multiplexed high-density electronic system, based on the VMM3a ASIC and the SRS system, is anticipated to mitigate significantly the limitations inherent in the current GET-based electronic system. This technological advancement is expected to enable the achievement of the spatial resolution demanded by the hypernuclei invariant mass spectroscopy experiment.

6.2.4. Displacement of Tracks Under Magnetic Fields

As illustrated in Figure 3.17, the magnetic field of GLAD exhibits non-uniformity with weak X and Y components inside the active volume. Furthermore, the potential non-horizontality of the TPC cannot be ruled out, as its alignment was solely based on a spirit level during the measurement preparation. The components of the magnetic field, not parallel to the drift field, introduce a transverse displacement of the drifting electrons upon reaching the anode. This process is described by the Langevin Equation, as formalized in Equations 2.32 to 2.34. Consequently, when in a magnetic field, the reconstructed laser tracks are expected to exhibit a deviation from their trajectories observed without a magnetic field. A comprehensive Monte Carlo simulation, which takes the field map of GLAD into consideration, is required to quantify accurately the influence of the magnetic field on the tracks.

Nevertheless, a preliminary study of the magnetic field's influence on track reconstruction was conducted using experimental data acquired under varying magnetic field strengths. The primary objective was to investigate the displacement of reconstructed laser tracks qualitatively. Data were collected at magnetic field strengths of approximately 0.3 T, 0.92 T, and 1.84 T. Due to a pronounced quenching effect, described in Section 6.2.2, at the highest field (1.84 T), no discernible peak was observed, limiting the investigation of the

magnetic field's influence at the maximum strength. For the remaining field strengths, the noise subtraction and de-multiplexing methods described in Sections 6.2.2 and 6.2.3 were applied to Peaks 1 and 2. The reconstructed laser track for Peak 1 under two different magnetic field conditions is presented in Figure 6.22. A collective shift of the laser tracks from Peak 1 towards lower X and Y indices in the pad plane with increasing magnetic field strength was observed. To better quantify the shift, fitting parameters and observables, namely the angles of the tracks and their intersection with the X-axis, are extracted from fitting and tabulated in Tables 6.2 and 6.3, respectively.

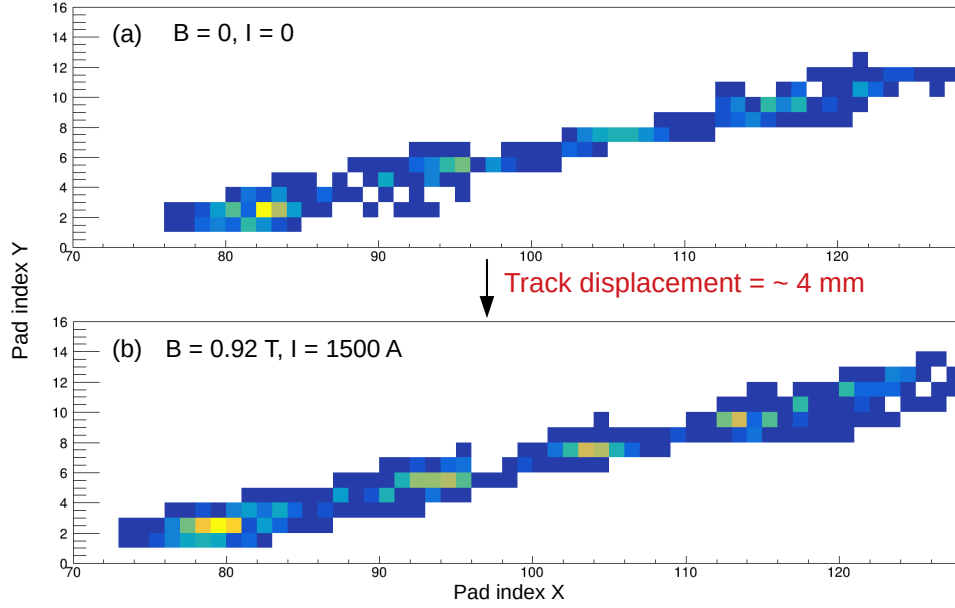


Figure 6.22.: Reconstructed tracks for laser Peak 1 from (a) Run 1 and (b) Run 3 of the first experiment.

The tables indicate a gradual displacement towards lower Intersection X points of several millimeters as magnetic field strength increases. These measured values deviate from the estimated electron drift displacement for both Peaks 1 and 2 using Equation 2.35 assuming a 1% magnetic field non-homogeneity (0, 1% B , B) (according to Figure 3.17), as summarized in Table 6.4. An explanation for this discrepancy lies in a possible tilt of the HYDRA-TPC at an angle of θ , as illustrated in Figure 6.23, as the horizontality of the TPC was solely verified using a spirit level as described in Section 6.2.1. Calculations demonstrate that for tilt angles ranging from 4.2° to 5.3° , the calculated displacement values exhibit excellent agreement with the measured values for both peaks and both magnetic field strengths, within the experimental uncertainty.

However, the comparison between the measurement and calculations cannot be further refined due to the large uncertainty. The underlying assumption of tracks' linearity under magnetic fields for reconstruction may not accurately reflect reality since the magnetic field inside the GLAD magnet is non-homogeneous. A comparative analysis accompanied by simulations, which include GLAD's field map, is required to assess the magnetic field's influence.

B [T]	Slope	Intercept [mm]	Angle α [$^\circ$]	Angle β [$^\circ$]	Intersection X [mm]
0	0.207 ± 0.007	-14.5 ± 0.7	11.67 ± 0.38	0.10 ± 0.81	139.8 ± 2.2
0.3	0.208 ± 0.004	-14.4 ± 0.4	11.74 ± 0.24	-0.53 ± 0.53	138.8 ± 1.4
0.92	0.209 ± 0.005	-14.1 ± 0.5	11.80 ± 0.28	-0.33 ± 0.55	135.3 ± 1.3

Table 6.2.: Fitting parameters and observables for Peak 1.

B [T]	Slope	Intercept [mm]	Angle α [°]	Angle β [°]	Intersection X [mm]
0	0.170 ± 0.003	9.8 ± 0.3	9.65 ± 0.19	-1.20 ± 0.35	114.7 ± 1.4
0.3	0.171 ± 0.003	9.7 ± 0.3	9.68 ± 0.19	-1.43 ± 0.38	114.0 ± 1.1
0.92	0.171 ± 0.007	9.3 ± 0.5	9.71 ± 0.42	-1.16 ± 0.50	109.3 ± 1.7

Table 6.3.: Fitting parameters and observables for Peak 2.

Peak	B [T]	Measurement [mm]	Calculation 1 [mm]	Calculation 2 [mm]	Calculation 3 [mm]
1	0.3	-1 ± 2.6	0.3	-0.9	-1.2
	0.92	-4.5 ± 2.6	0.7	-1.9	-2.6
2	0.3	-0.7 ± 1.8	0.7	-1.8	-2.5
	0.92	-5.4 ± 2.2	1.5	-4.1	-5.6

Table 6.4.: Comparison between the measured displacement and calculated values. B denotes the magnetic field in the GLAD magnet. Calculation 1 considers a positive magnetic field y component with a magnitude of 1% B . Calculation 2 considers a $\theta = 4.2^\circ$ tilting of the TPC in addition to Calculation 1. Calculation 3 considers a $\theta = 5.3^\circ$ tilting of the TPC in addition to Calculation 1.

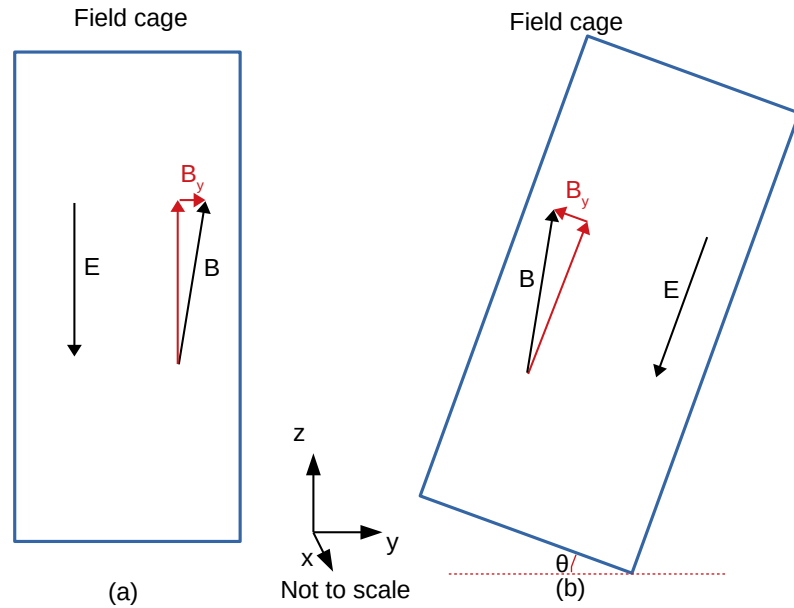


Figure 6.23.: Non-parallel component B_y of the magnetic field for (a) the horizontally placed TPC and (b) the tilted TPC at an angle θ .

6.2.5. Preliminary study for the Second Experiment

Experimental Setup

In the second experiment, measurements were conducted using a ^{12}C beam in addition to pure laser measurements. The laser beams during these measurements allowed for the characterization of space charge effects induced by the reaction fragments through variations of the laser tracks. The ^{12}C beam was directly extracted from SIS18, transported to Cave-C, and passed through the beam monitoring detectors

(LOS, ROLU, and two small fiber detectors as illustrated in Figure 3.1) before impinging on a 10 cm thick ^{12}C target. This target was positioned between the TPC and the magnet's entrance flange, as illustrated in Figure 6.24. The TPC and front-end electronics were installed in GLAD in a similar configuration to the first experiment, except for the distance of approximately 166 cm between the entrance flange and the nearest wires of the HYDRA-TPC. The target and the TPC were collimated with a laser source as shown in Figure 6.24.

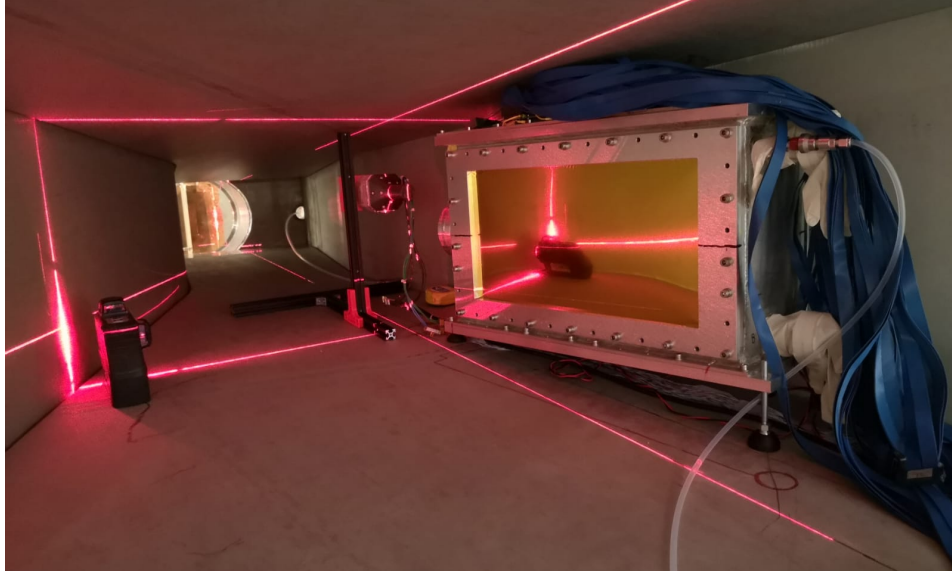


Figure 6.24.: Laser alignment of the ^{12}C target and the TPC for the second experiment.

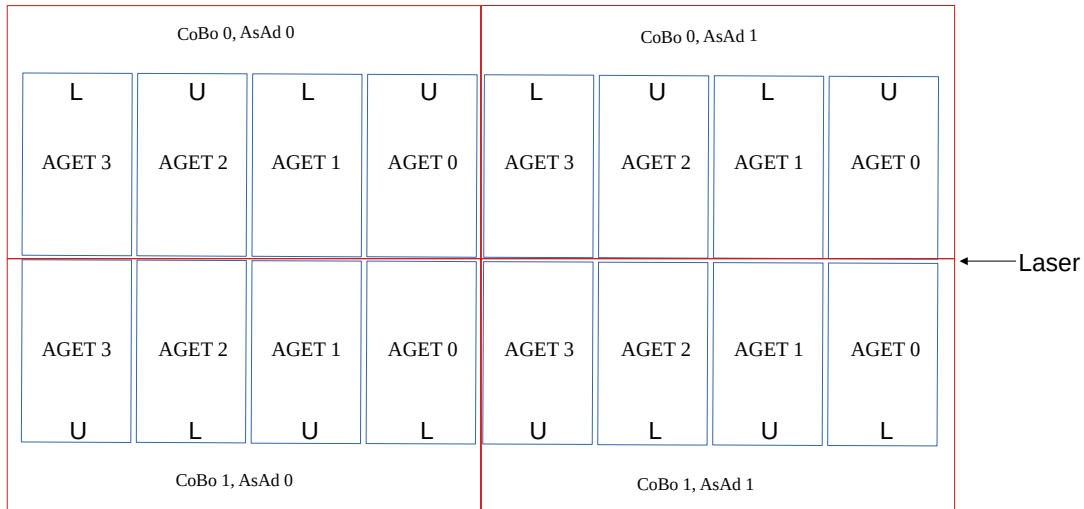


Figure 6.25.: Electronics mapping for the second experiment. 'U' and 'L' denote the socket connection for the multiplexing boards in Figure 3.20.

Furthermore, the scintillation wall was positioned between the TPC and the GLAD magnet wall. It was used to trigger the electronics on charged particles produced in $^{12}\text{C} + ^{12}\text{C}$ collisions that traversed the TPC and the wall. The UV laser source was positioned outside the magnet and emitted the laser beam directly into the TPC, which was different from the first experiment. The gas mixture and flow rate were maintained the same

as in the first experiment. The general high voltage settings were also identical to configurations in the first experiment while the general settings for the GET are as follows:

$$\begin{aligned}\text{Writing clock frequency} &= 12.5 \text{ MHz}, \\ \text{Trigger delay} &= 35 \text{ } \mu\text{s}, \\ \text{Peacking time} &= 232 \text{ ns}, \\ \text{Range} &= 120 \text{ fC}.\end{aligned}$$

The mapping for the second experiment is sketched in Figure 6.25.

Data Collection

Before receiving the ^{12}C beam during the second experiment, the current, I , in the GLAD magnet was gradually increased from 0 to 2900 A in consecutive steps. At each step, a measurement run was performed with the laser system as outlined in Table 6.5. The GET-based system was synchronized with the 20 Hz laser source, triggering simultaneous data acquisition from both CoBos.

Run Index	Current I [A]	Magnetic field B [T]	N
13	0	0	2312
14	1000	0.6	2393
15	2000	1.2	2114
16	2900	1.78	7334

Table 6.5.: Summary of data collected in the second experiment before accepting the ^{12}C beam. N represents the total number of events recorded by both CoBo 0 and CoBo 1.

Subsequently, measurements were performed with and without the laser and beam under the magnetic field of 2900 A. A comprehensive overview of the collected data, including the corresponding experimental conditions, is summarized in Appendix A.

Waveform analysis

During the second experiment, all five expected peaks were observed in the waveforms, as plotted in Figure 6.26, presenting an improvement over the first experiment. This is attributed to the direct laser injection into the TPC, superior to the use of the reflection mirror in the first experiment.

Peak amplitudes and times were determined by identifying the bins with the maximum ADC channel values. On CoBo 0, AsAd 1, AGET 2, for example, the mean values of peak time in units of time bins for the five peaks were calculated to be 198.3 ± 0.8 , 263.4 ± 0.6 , 339.8 ± 0.9 , 397.0 ± 1.5 and 448.1 ± 0.9 , respectively. As detailed in Section 6.1.2, these five peaks correspond to signals originating from Peak 0 of the GEM, Peak 4 of the cathode, and Peaks 1-3 of the three micro-mirror bundles. Therefore the drift velocity was calculated using the time difference between Peaks 0 and 4, expressed in units of time bins. The value was calculated to be $1.50 \pm 0.01 \text{ cm}/\mu\text{s}$. Furthermore, the distances between these peaks are calculated correspondingly: $78.2 \pm 1.2 \text{ mm}$ (GEM-to-first bundle), 91.7 ± 1.3 (first-to-second bundle), $68.7 \pm 2.1 \text{ mm}$ (second-to-third bundle) and $61.4 \pm 2.1 \text{ mm}$ (third bundle-to-Cathode).

Due to the possibly tilted orientation of laser tracks into the TPC, the distances between peaks in one AGET chip do not directly correspond to the actual distances between the micro-mirror bundles, the GEM, and the cathode. To determine the true distances, track reconstruction and back-extrapolation to the central axis of the mirrors, specified in Section 6.2.3, are necessary. For a simple estimation, the average peak time within each AGET chip is extracted and subsequently fitted with a linear function with the position of the central pads in each unit according to Figures 3.20 and 6.25. The calculated distances are summarized in

Table 6.6. For the first experiment, the distances were calculated utilizing the drift velocity determined in the second experiment and found to exceed the design values. This discrepancy may be attributed to variations in drift velocity between the two experiments, potentially arising from temperature fluctuations. In the second experiment, while the latter two distances demonstrate agreement within their respective uncertainties, the first two distances deviate by 5-6 mm from the design values. The disagreement is under further investigation. It may arise from a displacement of the first micro-mirror bundle.

Peak distance	First Experiment [mm]	Second Experiment [mm]	Design [mm]
0 - 1		79.4 ± 1.2	86.3 ± 2.5
1 - 2	94.0 ± 2.6	90.5 ± 1.3	85 ± 2.5
2 - 3	132.1 ± 3.4	68.7 ± 2.1	66 ± 2.5
3 - 4		61.4 ± 2.1	62.7 ± 2.5

Table 6.6.: Summary of distances between peaks for the first experiment, the second experiment, and from the design in Figure 6.2. '0 - 1' denotes the GEM-to-first bundle distance. '1 - 2' denotes the first-to-second bundle distance. '2 - 3' denotes the second-to-third bundle distance. '3 - 4' denotes the third bundle-to-cathode distance. Due to the absence of Peak 3 in the first experiment, the table displays the distance between Peaks 2 and 4.

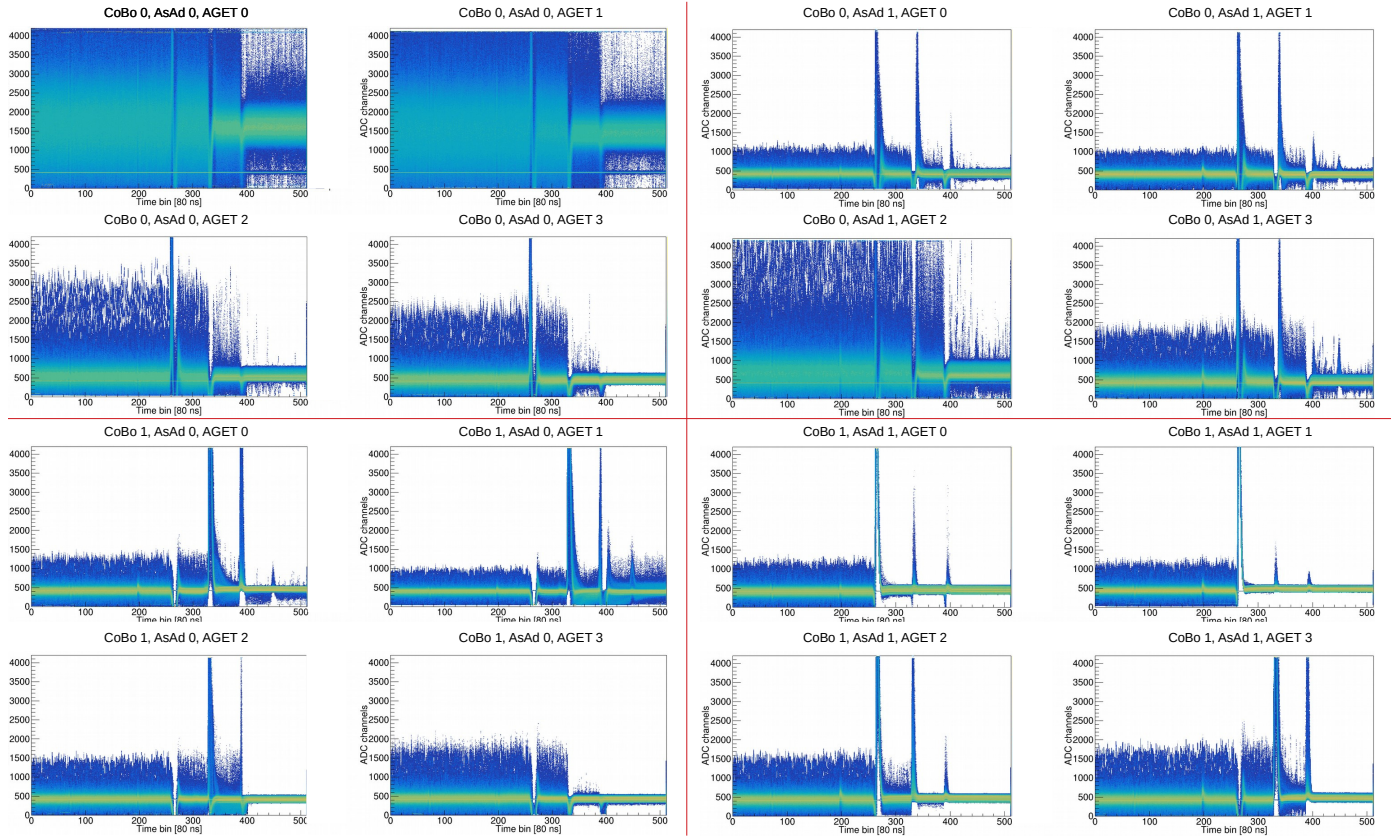


Figure 6.26.: Raw waveform for each AGET chip from all events of Run 13 in the second experiment.

A significantly higher noise level was observed compared to the first experiment, as illustrated in Figure 6.26. It was found to arise from poor cable connections for AGETs 0 and 1 of CoBo 0, AsAd 0, due to limited space

between the TPC and the GLAD ceiling. In other AGET chips, there are different levels of noise, although the source remains unidentified. For example, in Run 13, as illustrated in Figure 6.27, one event exhibited rapid fluctuations and decaying noise in most AGET chips. This was particularly pronounced in CoBo 0, AsAd 0, AGET 2, and CoBo 0, AsAd 1, AGET 2, where the noise significantly hindered the determination of peak time and amplitude. Consequently, the noise subtraction method in Section 6.2.2 is not directly applicable. Furthermore, as obvious in Figure 6.26, some signal peaks exhibit overflow, such as Peaks 1 and 2. The range of the GET configuration was not modified to capture the full peak during the measurement so that the low-amplitude peaks, such as Peaks 3 and 4, could be clearly distinguished from the significant noise. However, the presence of overflows and significant noise levels precluded the direct application of the analysis methods developed for the first experiment. An alternative analysis method is therefore required.

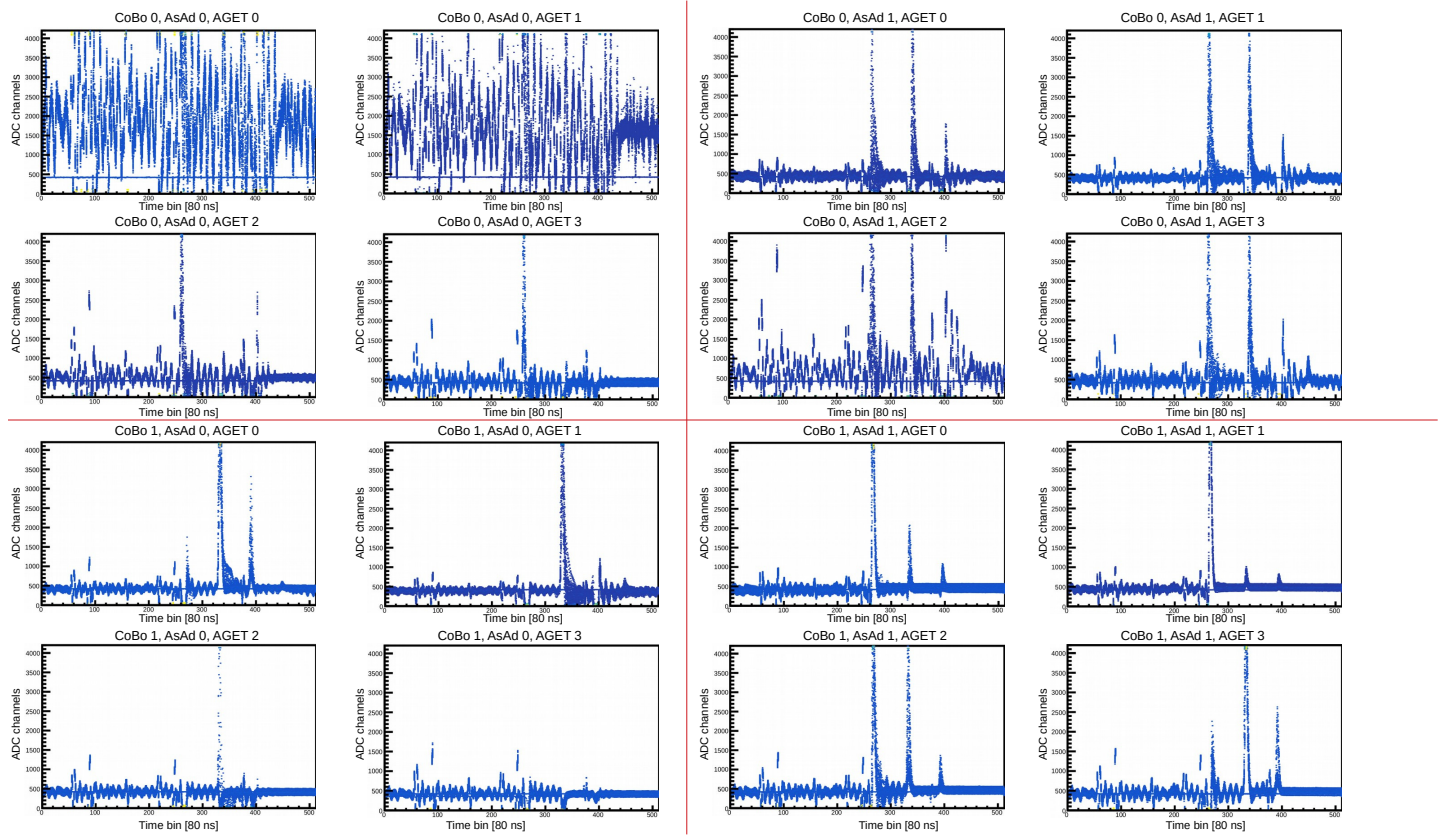


Figure 6.27.: Raw waveform for each AGET chip from one event of Run 13 in the second experiment.

6.2.6. Discussion

The HYDRA-TPC has been successfully commissioned with the laser system under various magnetic field conditions. A tracking algorithm was developed for the reconstruction of laser tracks, resulting in a spatial resolution better than 3 mm in the drift direction. However, the 200 μm resolution requirement in the pad plane is not perfectly met, primarily due to limitations imposed by the signal readout system with the intermediate multiplexing method. To address this, a system integrating VMM3 ASICs with the SRS readout is currently under development. This integration is expected to eliminate the need for multiplexing, thereby resolving the existing spatial resolution limitations.

The feasibility of tracking within the GLAD magnet using the HYDRA-TPC is demonstrated, alongside

an investigation into the influence of the magnetic field on electron drift. The measured electron drift displacements were found to be in good agreement with calculations based on the Langevin equation, considering a TPC tilt angle within the range of 4.2° to 5.3° . However, a more accurate assessment of the displacement under magnetic fields requires a direct comparison between experimental results and detailed simulations based on the field map of the GLAD magnet. Furthermore, the study of electron drift under the highest magnetic field strength (1.8 T) was precluded due to the signal quenching effect observed at this field strength. A complimentary laser measurement at this magnetic field strength is necessary.

To optimize the tracking performance of the TPC, further efforts are required in several aspects. Firstly, the underlying cause of the quenching effect on the laser signals remains elusive and requires further investigation. Secondly, discrepancies exist between the reconstructed and designed values for the angles of laser tracks and the position of the first micro-mirror bundle. Addressing these discrepancies demands precise measurements of these angles and distances. With precise knowledge of the laser track positions, a comprehensive evaluation of the drift field distortion can be conducted. Furthermore, the influence of the IBF and the resulting space charge effects on the HYDRA-TPC performance requires further assessment. This includes a more detailed analysis of the second experiment and future in-beam measurements where the full experimental setup is anticipated, incorporating both the TPC and the plastic wall.

7. Summary and Outlook

This work presents the development of a prototype Time Projection Chamber, the HYDRA-TPC for the HYDRA experiment within the R³B setup at GSI/FAIR. The TPC is a crucial component of a tracker, enabling high-precision invariant mass spectroscopy of light hypernuclei through the tracking of pions from the in-flight decay of the hypernuclei. To achieve a momentum resolution of 1%, spatial resolutions of 200 μm in the pad plane and 3 mm in the drift direction are required.

The prototype TPC, designed and constructed as described in Chapter 3, is a compact and lightweight detector with dimensions of $555 \times 229 \times 400 \text{ mm}^3$ and a mass of less than 35 kg. It is optimized for placement within the central region of GLAD's internal volume where the magnetic field is at maximum while keeping an estimated π^- detection efficiency of approximately 25%. It incorporates a two-layer wired drift field cage with a 300 mm drift distance and a $256 \times 88 \text{ mm}^2$ active area. The field cage encloses the drift region of the TPC together with a copper cathode and a GEM foil. A hybrid amplification stage comprising the GEM and a resistive Micromegas is employed, minimizing ion back-flow to less than 1%. The pad plane, with 128×44 pads, each pad $1.9 \times 1.9 \text{ mm}^2$, offers signal readout covering the entire active area. Furthermore, a laser system, incorporating three bundles of micro-mirrors, is embedded within the TPC to generate laser tracks inside the active region.

Two-dimensional simulations, detailed in Chapter 4, guided the field cage design, utilizing finite element and Monte Carlo methods to assess electric field uniformity. Results demonstrated that the aforementioned geometry achieves a maximum positional dispersion of approximately 210 μm and 230 μm at the active volume edges for the gas mixtures 90% Ar + 10% CO₂ and 90% Ar + 3% CF₄ + 2% iso-C₄H₁₀, respectively. In the central region, the dispersion falls within the desired 200 μm spatial resolution.

The gain performance of the TPC was characterized using an X-ray source in Chapter 5. Subsequently, laser measurements with the TPC were performed using 1024-channel GET-based electronics and intermediate multiplexing boards in Chapter 6. All major signal peaks from the micro-mirrors, GEM, and cathode were observed. By developing a track reconstruction method based on the multiplexing scheme, two laser tracks were successfully reconstructed within the active volume. The spatial resolution in the drift plane was found to be within 3 mm, however, the pad plane resolution did not meet the required 200 μm , mainly attributed to uncertainties arising from the multiplexing scheme as discussed in Section 6.2.3. Laser measurements were conducted under varying magnetic field conditions. The reconstructed laser tracks exhibited a displacement of several millimeters, which increased with the applied magnetic field strength, indicating a tilt of 4.2° to 5.3° for the TPC during the measurement. A detailed comparison between these measurements and simulations, based on the field map of the GLAD magnet, is yet pending.

While this work has conclusively demonstrated the functionality of the HYDRA-TPC and the feasibility of tracking under a non-homogeneous magnetic field, further efforts are required to enhance the TPC's performance and to understand several observations. To improve the spatial resolution in the pad plane, the implementation of the non-multiplexing continuous readout electronic system, comprising the VMM3 ASICs and SRS, is anticipated to mitigate the uncertainties introduced by the multiplexing scheme.

The position of the second signal peak deviated from the design value, and the orientation of reconstructed

laser tracks differed from pre-mounting measurements, potentially attributed to micro-mirror misplacement during installation. To address these issues, precise measurements of micro-mirror positions and laser track angles are crucial. Furthermore, with accurate knowledge of the laser track trajectories, a precise assessment of the drift field distortions can be conducted.

A comprehensive study of ion back-flow (IBF) and space charge effects within the TPC falls beyond the scope of this work, while minimizing them is critical for optimal TPC performance. To investigate these effects, a dedicated in-beam experiment is scheduled at CERN, utilizing a high-intensity particle beam. This experiment will employ the full experimental setup, incorporating the TPC, the plastic wall, and the electronics based on the VMM3 ASICs and SRS. This experiment will validate the functionality of the entire experimental setup under high readout rate conditions. The primary focus will be on the comprehensive characterization of the IBF and its resulting space charge distribution in the drift volume, which will be achieved through the analysis of laser track distortions.

Building upon the successful development of this prototype, a full-scale TPC is foreseen. With estimated dimensions of approximately $1000 \times 300 \times 200 \text{ mm}^3$, the full-scale TPC will incorporate all the features of the prototype while aiming for an improved π^- detection efficiency, ideally doubling the efficiency of the prototype. This detector will have broad applications beyond hypernuclear studies, including investigations into the nuclear shell structure, short-range correlations, fission studies, and flow measurements in heavy-ion collisions. Detailed designing and specific applications will be further refined through comprehensive simulations.

A. Data Overview for the February Measurement

Run Index	Type	Trigger rate	Beam intensity I_B	TD [μ s]	N	Others
17	L	LT = 20 Hz	0	35	10505	
18	B	PT = 4-5 kHz	5×10^4 *	31	6904	
19	B	PT = 4-5 kHz	5×10^4 *	43	6450	
20	B	PT = 4-5 kHz	5×10^4 *	50	6517	
21	B	PT = 4-5 kHz	5×10^4 *	35	467667	
22	L	LT = 20 Hz	0	35	90137	
23	L	LT = 20 Hz	0	40	3959	
24	L	LT = 20 Hz	0	40	613	range = 1 pC
25	B	PT = 4-6 kHz	5×10^4	0	10106	
26	B	PT = 4-6 kHz	5×10^4	5	4702	
27	B	PT = 4-6 kHz	5×10^4	10	8419	
28	B	PT = 1 kHz	4×10^3	0	1088	
29	B	PT = 1 kHz	4×10^3	15	2453	
30	B	PT = 1 kHz	4×10^3	20	974	
31	B	PT = 1 kHz	4×10^3	25	1158	
32	B	PT = 1 kHz	4×10^3	30	1013	
33	B	PT = 1 kHz	4×10^3	35	932	
34	L	LT = 20 Hz	0	35	89217	
35	L	LT = 20 Hz	0	35	1247	$U_{DLC} = 360$ V
36	L	LT = 20 Hz	0	35	11555	$U_{DLC} = 380$ V
37	LB	LT = 20 Hz	4.5×10^3	35	47712	
38	B	PT = 500 Hz	4.5×10^3	35	167588	I_B was unstable
39	B	PT = 500 Hz	4.5×10^3	30	33632	
40	LB	LT = 20 Hz	5×10^4 to 1×10^5	35	26239	I_B kept increasing
41	LB	LT = 20 Hz	1×10^5	35	15380	
42	B	PT = 1 kHz	1×10^5	30	53927	I_B was unstable
43	LB	LT = 20 Hz	5×10^5	35	34974	
44	B	PT = 1 kHz	5×10^5	30	168814	
45	LB	LT = 20 Hz	5×10^5	35	20018	

Table A.1.: Summary of data collected in February. ‘L’ represents the laser, ‘B’ the beam, ‘LB’ the simultaneous operation of laser and beam. ‘LT’ represents laser triggering, and ‘PT’ plastic wall triggering. The beam intensities in the table are expressed in units of particles per spill. ‘TD’ is short for trigger delay. Values marked with * were inferred from the counting rate of the LOS. The comment “ I was unstable” indicates that during the measurement the beam intensity decreased to half of the maximum value at some point.

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