

Chapter 2

Functional Principle of Micromegas

An introduction to Micromegas has already been given in Sect. 1.1. In the following chapter the functional principle as well as the main features of Micromegas detectors are discussed in detail. Readers interested in the internal setup of Micromegas detectors are referred to Sect. 2.1. The physics of ionization, drift and gas amplification processes are discussed in Sects. 2.2–2.4. In Sect. 2.5 the signal formation in Micromegas is described. Effects related to the transparency of the micro-mesh for electrons are discussed in Sect. 2.6. The streamer mechanism, that leads to the formation of discharges in Micromegas, is shortly described in Sect. 2.7.

2.1 Internal Setup

A Micromegas is a micro-pattern gaseous detector with strongly asymmetric drift and amplification regions. They are defined by three key components: a printed circuit board (PCB), carrying the readout structure, a fine micro-mesh, held by insulating pillars at a small distance of about 0.1 mm to the readout structure and a cathode, which closes the typically 6 mm wide drift region.

The readout structure typically consists of photo-lithographically etched copper strips with a width of 150 μm and a pitch of 250 μm , although more complex patterns like zig-zag-lines, pad or pixels are easily possible (Thibaud et al. 2014). Periodically spaced, insulating solder resist pillars with a diameter of 0.3 mm and a height of 128 μm define the width of the amplification region between anode and micro-mesh. Woven stainless steel meshes¹ have proven to be easy to handle during assembly, to sustain discharges and to permit a satisfactory energy and spatial resolution.

A several millimeter wide drift region is formed by the micro-mesh and a cathode, which can also directly close the gas filled region of the detector. Drift and amplification regions are filled with a suitable detector gas, usually a noble gas based mixture

¹E.g. 400 lines per inch, wire diameter 18 μm , wire periodicity 63.5 μm .

of e.g. Argon and Carbon Dioxide, Argon and Isobutane (Bay et al. 2002) or Neon, Ethane and Tetrafluoromethane gas (Bernet et al. 2005).

A comparison of different Micromegas types is given in Sects. 3.1 and 3.2. The internal setup of the Micromegas detectors, that have been developed in the context of this thesis, is described in detail in Sect. 3.4.

2.2 Interaction of Particles and Photons in the Detector

The underlying principle of each particle and photon detector is the interaction of the traversing particle or photon with the medium in the active volume of the detector. Since the electro-magnetic cross sections dominate by orders of magnitude over the weak, strong or even gravitational cross sections, five different interaction mechanisms of charged particles are used in gas detectors: Ionization, excitation, production of Čerenkov radiation, bremsstrahlung and production of transition radiation.

The detection mechanism in Micromegas is the ionization of the detector gas by traversing particles or photons. Photons with an energy above several 100 eV deposit energy indirectly over production of charged particles.

Micromegas, as most gas detectors, are usually insensitive to neutral particles such as neutrons or neutral mesons. This improves the detector tracking performance in high-rate neutron background environments such as the ATLAS Small Wheel region (ATLAS Collaboration 2013). If desired, the efficiency to neutrons can be increased by the use of special converter media with high neutron capture cross sections, where the neutron produces a charged particle which is subsequently detected. Thermal neutrons are efficiently detected in ^3He -filled detectors via the $^3\text{He}(n,p)^3\text{H}+764\text{ keV}$ reaction. In the same way the $^{10}\text{B}(n,\alpha)^7\text{Li}+2.792\text{ MeV}$ reaction is used, where the boron can either be gaseous boron trifluoride or solid elementary boron.

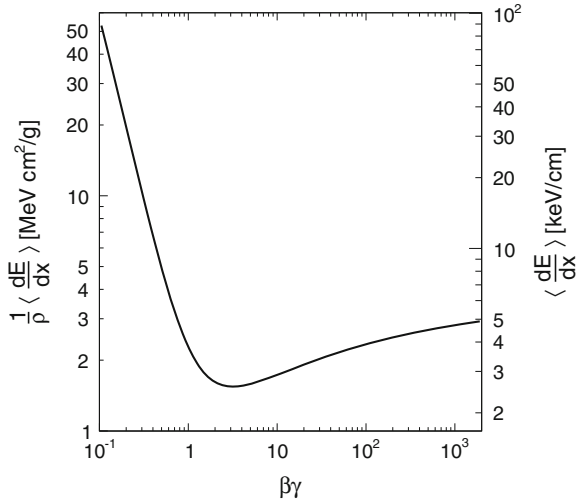
2.2.1 Charged Particles

The interaction of charged particles with matter via ionization, excitation and Čerenkov radiation by exchange of (virtual) photons is described in the Photo Absorption Ionization model (Allison and Cobb 1980). It leads to the Bethe-Bloch formula, which describes the mean energy loss per unit length²

$$\left\langle \frac{dE}{dx} \right\rangle = -4\pi r_e^2 m_e c^2 \rho N_A \frac{Zz^2}{A\beta^2} \left(\frac{1}{2} \ln \frac{2m_e c^2 \beta^2 \gamma^2 T_{\max}}{I^2} - \beta^2 - \frac{\delta}{2} - \frac{C}{Z} \right), \quad (2.1)$$

² r_e : classical electron radius, m_e : electron mass, ρ : density of the target material, N_A : Avogadro constant, Z : atomic number and A : atomic mass in g/mol of target material, z : charge of incident particle, $\beta = v/c$: velocity of incident particle, $\gamma = 1/\sqrt{1-\beta^2}$: Lorentz factor, $T_{\max} \approx 2m_e c^2 \beta^2 \gamma^2$: maximum kinetic energy transferable to an electron in an elastic collision, I : mean excitation energy.

Fig. 2.1 Mean energy loss of a singly charged particle in an Ar:CO₂ 93:7 gas mixture at 20°C and 1013 mbar in the range $0.1 \leq \beta\gamma \leq 2000$. The Barkas-Berger shell corrections for small $\beta\gamma$ and radiative corrections i.e. bremsstrahlung for large $\beta\gamma$ have been neglected



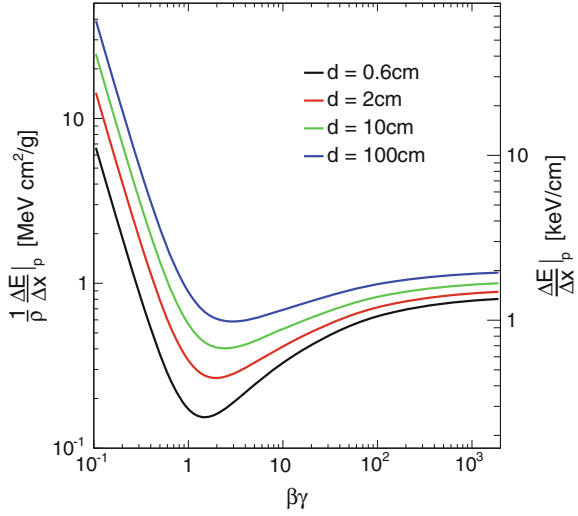
with the density and atomic structure corrections in the Sternheimer-Peierls and Barkas-Berger parametrization respectively. Tabulated values for the density correction δ and the mean excitation energy I for various materials can be found in Sternheimer et al. (1984). For compounds or mixtures the mean energy loss can be calculated from the weight fraction w_j of the j th component

$$\frac{1}{\rho} \left\langle \frac{dE}{dx} \right\rangle = \sum_j \frac{w_j}{\rho_j} \left\langle \frac{dE}{dx} \right\rangle_j. \quad (2.2)$$

The mean energy loss of a singly charged particle in an Ar:CO₂ 93:7 vol% gas mixture at normal temperature and pressure, calculated from Eqs. (2.1) and (2.2), is shown in Fig. 2.1. Its universal shape is best visible by plotting it as a function of $\beta\gamma$. The mean energy loss decreases approximately like $1/\beta^2$ for $\beta\gamma \lesssim 1$ and has a minimum with only small variations for $2 \lesssim \beta\gamma \lesssim 8$. Particles with a velocity in this region are called minimum ionizing. For almost all materials, the mean energy loss of minimum ionizing particles is on the order of 2 MeV cm²/g. For increasing $\beta\gamma$ the mean energy loss rises like $\ln(\beta^2\gamma^2)$ and reaches the so called Fermi plateau for $\beta\gamma \gtrsim 500$.

It has been argued by Bichsel (2006), that the mean energy loss $\langle dE/dx \rangle$, described by Eq. (2.1), is not directly accessible by measurements in thin detectors, since large energy transfers to target electrons strongly influence the measured values. In particle identification detectors, that use the repeated measurement of energy loss, methods like the so called truncated mean have been applied to estimate the mean energy loss, where the set of highest measured values is discarded. It is discussed by Beringer et al. (Particle Data Group) (2012) to rather use the most probable energy loss $\Delta E/\Delta x|_p$,

Fig. 2.2 Most probable energy loss of a singly charged particle in an Ar:CO₂ 93:7 gas mixture at 20 °C and 1013 mbar for gas volumes with different thickness d



formulated in Landau's theoretical description of ionization in gas detectors. For charged particles in a thin detector with thickness d we get

$$\left. \frac{\Delta E}{\Delta x} \right|_p = 4\pi r_e^2 m_e c^2 N_A \frac{Zz^2 \rho}{A\beta^2} \left(\frac{1}{2} \ln \frac{2m_e c^2 \beta^2 \gamma^2 \xi}{I^2} + \frac{j}{2} - \frac{\beta^2}{2} - \frac{\delta}{2} \right), \quad (2.3)$$

where $\xi = 2\pi r_e^2 m_e c^2 N_A \frac{Zz^2 \rho d}{A\beta^2}$ and $j = 0.200$ (Bichsel 1988; Grupen and Shwartz 2008). It should be noted, that the most probable energy loss $\Delta E/\Delta x|_p$ depends, unlike the mean energy loss $\langle dE/dx \rangle$, on the thickness d or ρd of the sensitive volume. For increasing thickness, it approaches the mean energy loss, described by the Bethe-Bloch formula. In Fig. 2.2 the most probable energy loss for an Ar:CO₂ 93:7 vol% gas mixture, calculated from Eqs. (2.2) and (2.3), is shown.

Although the straggling of the most probable energy loss in thin gas detectors is considerably larger than predicted by the Landau-Vavilov-Bichsel theory i.e. the measured pulse height spectra are not adequately described by a Landau function, the maximum of the distribution is correctly described by Eq. (2.3). In order to extract the measured most probable pulse height from the distribution of pulse heights, the distribution is fitted with a Landau function, convoluted with a Gaussian. The straggling can thus be absorbed into the Gaussian function.

The ionization yield i.e. the total number of produced electron-ion pairs can then be calculated from the mean or most probable energy loss ΔE in the active area of the gas detector

$$n_{\text{tot}} = \frac{\Delta E}{W_I}, \quad (2.4)$$

Table 2.1 Properties of typical elementary and molecular detector gases at 20 °C and 1013 mbar (Beringer et al. (Particle Data Group) 2012; Sternheimer et al. 1984; Grupen and Shwartz 2008)

Gas	Z	Z/A (g/mol)	ρ (g/cm ³)	I (eV)	W_I (eV)	n_p (cm ⁻¹)	n_t (cm ⁻¹)
He	2	0.49967	1.79×10^{-4}	41.8	41.3	3.5	8
Ne	10	0.49556	8.3851×10^{-4}	137.0	37	13	40
Ar	18	0.45059	1.6620×10^{-3}	188.0	26	25	97
Kr	36	0.42959	3.4783×10^{-3}	352.0	24	22	192
Xe	54	0.41130	5.4854×10^{-3}	482.0	22	41	312
CO ₂	–	0.49989	1.8421×10^{-3}	85.0	34	35	100
CH ₄	–	0.62334	6.6715×10^{-4}	41.7	30	28	54
C ₂ H ₆	–	0.59861	1.2532×10^{-3}	45.4	26	48	112

Z atomic number, only for elementary gases, Z/A (average) atomic number/atomic mass, ρ density, I mean excitation energy, W_I average energy per created electron-ion pair, n_p and n_t mean number of primary and total number of electron-ion pairs created by a minimum ionizing particle of unity charge

where W_I is the average energy, necessary for the production of one electron-ion pair. W_I is for most materials more than twice as high as the ionization potential, since energy is lost due to excitation of the gas atoms or molecules. It is, similar to the energy loss, Bragg additive, such that it can be calculated for mixtures from a relation similar to Eq. (2.2). Higher energetic secondary particles or photons, leaving the detector active volume without depositing their full energy, can complicate the relation between the calculated energy loss and the detectable ionization yield.

A compilation of typical detector gases and their properties can be found in Table 2.1.

Čerenkov radiation is emitted by relativistic charged particles in media with refractive index n , when their velocity exceeds the speed of light in the medium

$$\beta \geq \frac{1}{n} . \quad (2.5)$$

The emission can be understood as the coherent sum of the electric fields of molecular dipoles, created by the polarization of the medium by the traversing particle. As the particle is faster than the speed of light, the sum is non-zero (Čerenkov 1937).

Čerenkov radiation is emitted under a specific angle

$$\cos \vartheta_{\check{c}} = \frac{1}{\beta n} \quad (2.6)$$

to the particle direction of flight. The number of emitted photons can be calculated from the Frank-Tamm formula (Grupen and Shwartz 2008) and is for the visible range between $\lambda = 400$ and 700 nm given by

$$\frac{dN}{dx} = 490 \sin^2 \vartheta_{\check{c}} \text{ cm}^{-1} . \quad (2.7)$$

The emission of a high-energetic photon, due to the deceleration of a particle with an energy E above the so called critical energy in the proximity of a nucleus, is called bremsstrahlung. Although bremsstrahlung in gases is encountered in e.g. the interaction of electrons and positron with residual gas molecules and ions in storage rings and imposes additional radiation protection measures (Ipe and Fassò 1994), the cross sections in gases are too low to be relevant in gas detectors. The radiation length for Argon at normal temperature and pressure is with $X_0 = 120$ m much larger than typical detector dimensions.

Transition radiation with an energy of several keV is emitted by a highly relativistic particle traversing a boundary with alternating refractive indices. This can be understood as the emission of radiation from the variable dipole, formed by the incoming particle with its mirror charge. In practice, several layers consisting of e.g. carbon or Mylar are stacked to increase the photon yield (Beringer et al. (Particle Data Group) 2012). Due to interference effects, the transition radiation yield saturates at certain values. Transition radiation, produced in a radiator between the straw tubes in the Xenon filled ATLAS Transition Radiation Tracker, is used for electron/pion separation (Sect. 1.2.1). There are ideas to use Xenon filled Micromegas detectors, covered with a suitable radiator, in the same way.

2.2.2 Photons

Unlike the interaction of particles with matter, the interaction of photons, collimated to a narrow beam, is mostly a binary process. A photon either interacts with matter and is lost or re-emitted at lower energy or scattered out of the beam or it does not interact.

The intensity of a mono-energetic narrow photon beam with initial intensity I_0 after traversing a material with thickness d is given by

$$I(x) = I_0 \exp(-\mu d) = I_0 \exp\left(-\frac{\mu}{\rho} x\right), \quad (2.8)$$

where ρ is the density, $x := d\rho$ the mass thickness and μ/ρ the mass attenuation coefficient of the material. The energy dependent mass attenuation coefficient is related to the interaction cross section $\sigma(E)$ via

$$\frac{\mu(E)}{\rho} = \frac{\sigma(E)}{m}, \quad (2.9)$$

where m is the atomic mass of the absorber material.

For mixtures and compounds, an effective mass attenuation coefficient can be calculated

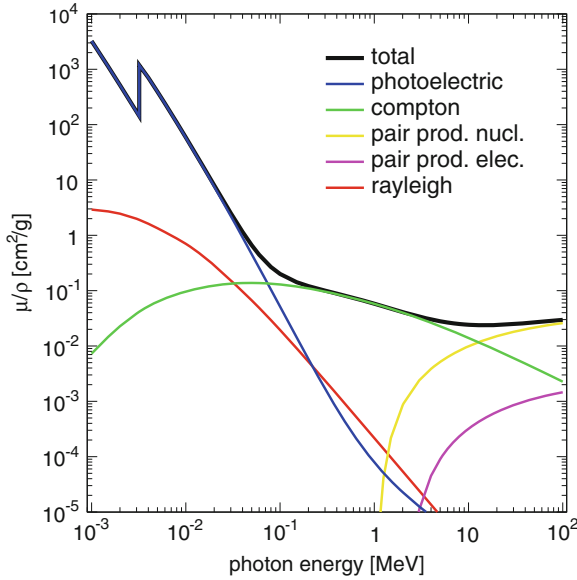


Fig. 2.3 Mass attenuation coefficient as a function of the photon energy for an Ar:CO₂ 93:7 vol% gas mixture. The data are taken from Deslattes et al. (2010). Photoelectric effect, Compton scattering and nuclear pair production dominate the mass attenuation coefficient for low, medium and high photon energies. Coherent Rayleigh scattering of photons on electrons, i.e. without excitation or ionization of the material and pair production in the electron field, add small contributions for low and high photon energies respectively

$$\frac{\mu}{\rho_{\text{eff}}} = \sum_i w_i \frac{\mu}{\rho_i} , \quad (2.10)$$

with the weight fraction w_i of the i th component with coefficients μ/ρ_i .

Three main processes dominate the underlying photon interaction with matter: Photoelectric effect, Compton scattering and electron-positron pair production (Deslattes et al. 2010), see Fig. 2.3.

Photoelectric effect, the absorption of a photon with energy E_γ by an atom, accompanied by the release of a shell electron, dominates the interaction cross section for low photon energies up to several 100 keV. The emitted electron energy is given by

$$E_{e^-} = E_\gamma - E_b , \quad (2.11)$$

where E_b is the binding energy of the electron. If the emitted electron comes from a lower shell, the atom is left in an excited state and will typically deexcite by emission of low energy photons with total energy $\sim E_b$. If a photon interacts via photo-effect in the active volume of a gas detector, and the active region is thick enough to also detect the photons from the atom deexcitation, the whole incident photon energy is deposited in the detector.

The photoelectric cross section strongly depends on the atomic number Z of the material

$$\sigma_{\text{photo}} \propto Z^5, \quad (2.12)$$

such that high- Z materials like Xenon can be used in gaseous photon detectors to considerably increase the detection efficiency (Kleinknecht 1992).

Compton effect can be interpreted as elastic scattering of photons on quasi free electrons. From energy and momentum conservation, the energy of the scattered photon can be calculated as

$$E'_\gamma = \frac{E_\gamma}{1 + (1 - \cos \vartheta) E_\gamma / m_e c^2}. \quad (2.13)$$

Thus the scattered electrons possess a maximum kinetic energy

$$T'_{e,\text{max}} = E_\gamma \frac{2E_\gamma / m_e c^2}{1 + 2E_\gamma / m_e c^2}, \quad (2.14)$$

which leads to the so called Compton edge in the energy spectrum.

Photons with $E_\gamma \geq 2m_e c^2$ can create an electron-positron pair in the proximity of a nucleus. This process is called nuclear pair production. The nucleus is necessary for simultaneous energy and momentum conservation. The analogous process in the proximity of an electron is accordingly called electron pair production.

2.2.3 Multiple Scattering of Charged Particles

Due to the numerous but small interactions of charged particles with the detector components and the detector gas, the particle direction of flight is altered. This effect is called multiple scattering and is typically unwanted in particle detectors. It leads to an increased divergence of a particle beam with initially parallel particle momenta.

The width θ_0 of the Gaussian shaped angular distribution, resulting from multiple scattering of a narrow beam of particles with parallel momenta p and charge number z in a material of thickness x and radiation length X_0 is given by Beringer et al. (Particle Data Group) (2012)

$$\theta_0 = \frac{13.6 \text{ MeV}}{\beta c p} z \sqrt{\frac{x}{X_0}} \left(1 + 0.038 \ln \left(\frac{x}{X_0} \right) \right). \quad (2.15)$$

The angular scattering can be translated into a transverse broadening of the particle beam. Its width after a distance d to the scattering object is approximated by

$$\Delta y = \frac{\theta_0}{\sqrt{3}} d. \quad (2.16)$$

For a composite scatterer, consisting of layers of materials with different radiation lengths X_i and densities ρ_i and thickness d_i , the resulting radiation length can be calculated:

$$\frac{1}{X_0} = \sum \frac{w_i}{X_i}, \quad (2.17)$$

where the weight fraction $w_i = d_i \rho_i / \sum d_j \rho_j$ (Gupta 2013).

Multiple scattering distorts the measured spatial resolution in Micromegas detectors, as often straight tracks are assumed as prerequisite. For many applications of Micromegas in low and medium energy particle tracking, the detector material budget is reduced wherever possible, to decrease the degradation of the particle beam.

2.3 Drift of Electrons and Ions in Gases

The drift of electrons and charged ions in gases is governed by electric and magnetic fields $\vec{E}$ and $\vec{B}$. The drift velocity for electrons is given by

$$\vec{v}_d = \frac{e}{m_e} \frac{\tau}{1 + \omega^2 \tau^2} \left(\vec{E} + \frac{\omega \tau}{B} (\vec{E} \times \vec{B}) + \frac{\omega^2 \tau^2}{B^2} (\vec{E} \cdot \vec{B}) \vec{B} \right), \quad (2.18)$$

where e and m_e are electron charge and mass, respectively, $\omega = eB/m_e$ is the Larmor frequency and τ the mean time between two collisions with gas atoms (Beringer et al. (Particle Data Group) 2012).

Three properties of the electron drift can be deduced from Eq. (2.18).

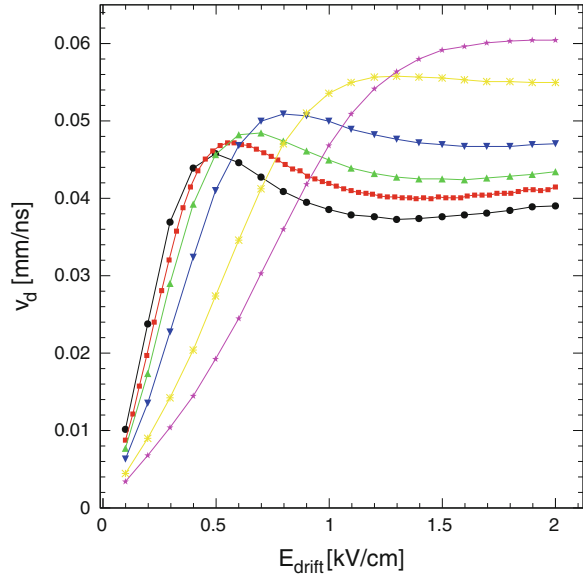
First, in the absence of a magnetic field, the electrons follow on average the electric field lines. Electron diffusion transverse and parallel to the electric field, not considered in the equation, leads to a broadening of an initially point-shaped charge distribution.

Second, for perpendicular electric and magnetic fields, the $\vec{E} \times \vec{B}$ -term leads to a drift component perpendicular to the fields, such that the net drift velocity vector spans the so called Lorentz angle $\alpha_L \approx \arctan(\omega \tau)$ with the electric field vector. This effect leads to a systematic shift of hit positions with perpendicularly incident particles and compresses or disperses the detected charge distribution for inclined tracks in Micromegas used in a magnetic field.

Third, for non perpendicular electric and magnetic field, there is a drift component into the direction of the magnetic field. For large magnetic fields, i.e. $\omega \tau \gg 1$ the net drift vector may even point into the same direction as $\vec{B}$. Diffusion of electrons perpendicular to the magnetic field is strongly suppressed, as the field forces the electrons on a helix around the magnetic field lines. In Time-Projection-Chambers, this is used to improve the spatial resolution, which would suffer from large electron diffusion.

In the following, the magnetic field is assumed to be zero.

Fig. 2.4 Electron drift velocity as a function of the electric drift field for various Ar:CO₂ gas mixtures at 20 °C and 1013 mbar, computed with MAGBOLTZ (Biagi 1999). The color coding is explained in Fig. 2.5



Due to their small mass, electrons gain sufficient energy in the electric field such that their de-Broglie-wavelength is of the same order as typical atomic dimensions, which leads to a strong energy and electric field dependence of their interaction with the gas.

In practice, electron drift velocities as well as diffusion coefficients as a function of the electric field are calculated with MAGBOLTZ (Biagi 1999), see Figs. 2.4 and 2.5. Usually mixtures of gases are used in order to optimize the detector performance. In this thesis gas mixtures are defined by the volumetric mixing ratios of the constituent gases.

The drift of ions can be more easily described by

$$v_{d,\text{ion}} = \mu_{\text{ion}} E , \quad (2.19)$$

where μ_{ion} is the gas dependent ion mobility, which is constant over a wide electric field range and not strongly influenced by the admixture of molecular gases. Mobilities for several ions in gases are compiled in Table 2.2.

The drift of ions from gas amplification from the anode strips towards the mesh defines, together with the electron drift time in the drift region, the intrinsic length of signals in Micromegas. In order to further improve the high-rate capabilities by decreasing the signal length, the use of gases lighter than Argon is necessary. This has furthermore an advantageous influence on the discharge probability, as a quick drain of space charge reduces the streamer formation probability, see Sect. 2.7.

Fig. 2.5 Transverse diffusion coefficient as a function of the electric drift field for various Ar:CO₂ gas mixtures at 20 °C and 1013 mbar, computed with MAGBOLTZ. The volumetric mixing ratios are stated in the legend. From variation of the Monte Carlo integration parameters, used in MAGBOLTZ, an accuracy for the shown coefficients of the order 5 % is approximated

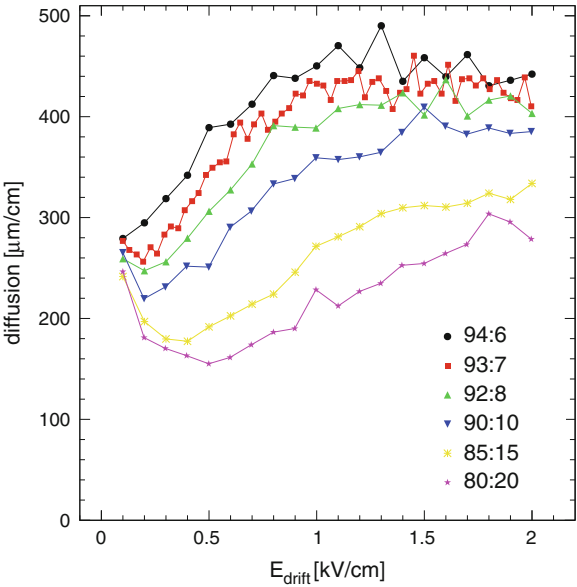


Table 2.2 Mobility of ions in gases at 20 °C and 1013 mbar (MacDaniel and Mason 1973)

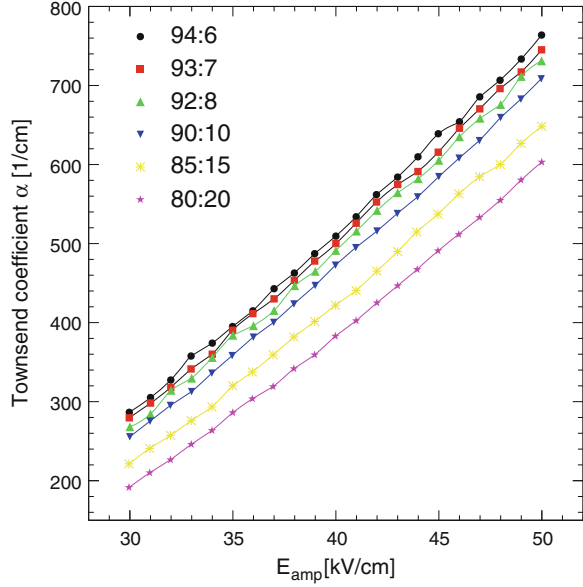
Gas	Ion	$\mu_{ion}[cm^2/(Vs)]$
He	He ⁺	10.4
Ne	Ne ⁺	4.7
Ar	Ar ⁺	1.54
Ar:CH ₄	CH ₄ ⁺	1.87
Ar:CO ₂	CO ₂ ⁺	1.72
CH ₄	CH ₄ ⁺	2.26
CO ₂	CO ₂ ⁺	1.09

2.4 Gas Amplification

The ionization charge, created by minimum ionizing particles in gas detectors, is on the order of only 100 e/cm, Table 2.1. Thus in order to reliably detect the passage of singly charged particles or photons, a process of charge amplification is needed. In the following, charge amplification in gas detectors by electron impact ionization, secondary emission and Penning transfer is discussed.

Free electrons from e.g. ionization processes, accelerated by an appropriate electric field, can gain a sufficient kinetic energy between two collisions with gas atoms to further ionize the gas atoms or molecules by impact ionization.

Fig. 2.6 First Townsend coefficient as a function of the amplification field for different Ar:CO₂ gas mixtures at 20 °C and 1013 mbar, calculated with the program MAGBOLTZ (Biagi 1999). Note that Penning transfer is not incorporated in the shown values



Assuming n_0 electrons entering the high electric field region between mesh and anode strips of a Micromegas detector, the total number of electrons after a path length x is given by

$$n(x) = n_0 \exp(\alpha x) . \quad (2.20)$$

The first Townsend coefficient α is a function of the electric field E , the gas density and the gas type. Calculated Townsend coefficients for several gas mixtures are shown in Fig. 2.6.

The so called gas gain is the factor between the initial and the final charge and is defined as

$$G = \exp(\alpha d) , \quad (2.21)$$

where d is the total width of the amplification region.

The density of the detector gas considerably influences the gas gain. For Argon based mixtures, the relation between gas density n , temperature T and pressure p is given by the ideal gas law

$$n = \frac{p}{k_B T} , \quad (2.22)$$

where k_B is the Boltzmann constant.

In measurements with a standard Micromegas under controlled gas temperature and pressure conditions by Lippert (2012), supervised by the author, the gas density dependence of the gas gain (Eq. 2.21) has been investigated. The first Townsend coefficient can be parametrized with an empirical relation by Townsend (1910)

$$\alpha = A_0 n \exp\left(-\frac{B_0 n}{E_{\text{amp}}}\right), \quad (2.23)$$

where $A_0 = k_B A$ and $B_0 = k_B B$ are gas dependent parameters. For an Ar:CO₂ 93:7 vol% gas mixture, mean parameters of $A = (111.2 \pm 0.6) \text{ K}/(\text{bar} \mu\text{m})$ and $B = (2196 \pm 7) \text{ K V}/(\text{bar} \mu\text{m})$ have been found (Lippert 2012).

If the electric field in the amplification region is sufficiently large, impact ionization by positive ions or photo-effect by photons, produced in gas de-excitation, can lead to secondary emission of electrons from the mesh. This is in Micromegas typically only relevant, if the amplification field is considerably enhanced during discharge formation.

Şahin et al. (2010) state, that the observed gas gains in several gas mixtures are significantly higher than the value predicted by the Townsend theory, due to Penning like transfers i.e. the transformation of excitation energy of one to ionization of the other gas constituent. Penning and Penning-like transfers are possible, if constituent A can be excited to energy levels higher than the ionization threshold of B . Several transfer processes are possible for binary mixtures with constituents A and B , see Kuger (2013).

Deexcitation processes of A^* , competing with Penning transfers, are inelastic excitation of B to energy levels below the ionization energy e.g. excitation of rotational or vibrational energy levels, or radiative decays of A^* and excimers via photons, leaving the active detector area.

The contribution of Penning transfers to the overall gas gain can be quantified by the transfer rate r , describing the fraction of excited atoms or molecule of species A in states, eligible for direct transfer, that lead to an ionization of B . Şahin et al. (2010) claim a transfer rate for an Ar:CO₂ 93:7 vol% of (0.42 ± 0.03) . Upon comparing gas gain measurements by Kuger (2013) and the author, with a GARFIELD++ simulation (Veenhof 2010) incorporating Penning transfers with the stated transfer rate, a good agreement between measured and calculated gas gains is shown. Measurements conducted by Moll (2013) with strongly ionizing alpha particles in a standard, non-floating-strip Micromegas, confirm the correct description of gas amplification by the GARFIELD++ program, with the Penning transfer rate, stated above.

Typical gas gains between 600 and 10^4 are used in Micromegas.

The measurement of absolute gas gains in Micromegas is possible with two different methods:

1. For low rate irradiation, a measurement of the pulse height of charge signals on the anode with charge sensitive preamplifiers allows for a determination of the gas gain. Due to the considerable fluctuations in energy loss of charged particles in gas detectors and the dependence of the most probable energy loss on the

exact drift gap width (Sect. 2.2.1), energy deposition by low energy X-rays yields more reliable results. Low energy photons e.g. 5.9 keV X-rays, produced in the β -decay of ^{55}Fe , interact via photo-effect with the detector gas (Sect. 2.2.2). The produced photoelectron deposits its kinetic energy within 0.5 mm and produces on average an ionization charge of $5.9 \text{ keV} / 26.56 \text{ eV} = 222 e = q_0$. Considering the electron mesh transparency $t(E_{\text{drift}})$, the charge-to-voltage conversion factor of the preamplifier $c_{q \rightarrow U}$ and a detector capacitance dependent factor c_{cap} (see Bortfeldt 2010, Chap. 8 for a discussion), the pressure and temperature dependent gas gain $G(E_{\text{amp}}, T, p)$ can be calculated from the measured pulse height $ph(E_{\text{amp}}, E_{\text{drift}}, T, p)$ from

$$G(E_{\text{amp}}, T, p) = \frac{ph(E_{\text{amp}}, E_{\text{drift}}, T, p)}{q_0 t(E_{\text{drift}}) c_{\text{cap}} c_{q \rightarrow U}}. \quad (2.24)$$

This method relies on the accurate knowledge of the preamplifier specific factor $c_{q \rightarrow U}$ and the detector and readout electronics specific factor c_{cap} .

2. At high-rate irradiation with charged particles or low energy photons, a constant current of ionization charge is produced. Measuring the resulting current between anode and mesh $I(f, E_{\text{drift}}, E_{\text{amp}}, T, p)$ allows for an alternative gas gain determination. The gas gain

$$G(E_{\text{amp}}, T, p) = \frac{I(f, E_{\text{drift}}, E_{\text{amp}}, T, p)}{q(E_{\text{drift}}) f} \quad (2.25)$$

depends on the particle rate f and the charge, reaching the amplification region $q(E_{\text{drift}}) = q_0 t(E_{\text{drift}})$, which is given by the ionization yield q_0 and the drift field dependent electron mesh transparency $t(E_{\text{drift}})$. This method is insensitive to preamplifier related calibration, such as the charge-to-voltage conversion and detector capacitance correction factor as well as preamplifier shaping time. For low energy particles and long particle tracks in the detector, the predicted mean energy loss can be used to calculate the ionization yield. Care must be taken, that high-rate effects such as saturation, anode charge-up, ion space-charge in the amplification region or ion backflow corrections are taken into account.

Absolute gas gain factors are inherently difficult to measure with Micromegas detectors: The pulse height of signals, measured with charge sensitive preamplifiers at the anode, is strongly influenced by the internal capacitances within the detector and less strongly by the charge-to-voltage conversion factors of the preamplifiers (Bortfeldt 2010; Bortfeldt et al. 2012). In the alternative continuous current mode, where the constant current between mesh and anode strips, caused by continuous ionization of a high-rate particle or photon beam is measured, these calibration effects can be avoided. They are unfortunately replaced by distortions of the gas gain due to the charge-up, ion backdrift and other high-rate effects.

In this thesis, the absolute gas gain has been determined e.g. with 20 MeV protons (Sect. 6.1). But whenever possible, the relative gas gain as a function of the amplification field is used for a discussion of the results.

2.5 Signal Formation

The formation of signals in Micromegas has been discussed in detail in Bortfeldt (2010) and Bortfeldt et al. (2012) and is shortly summarized here. The description is inspired by a discussion of signal formation in parallel plate avalanche chambers by Mathieson and Smith (1988).

Charge signals are created by ionization processes due to traversing particles or photons followed by gas amplification and subsequent drift of ions from the anode to the mesh.

Consider a localized energy deposition in the several millimeters wide drift region e.g. by a X-ray photon with 5.9 keV, resulting in a cloud of electrons and positive ions, with a diameter on the order of 0.5 mm. Electrons and ions are separated by the electric drift field, the electrons drift with $v_d \sim 0.04$ mm/ns towards the high-field region between mesh and anode (Sect. 2.3). Depending on the dimension of the drift field, a certain fraction of electrons can enter the amplification region, the rest is lost on the micro-mesh (Sect. 2.6). Upon entering the typically 128 μ m wide region between mesh and anode, the electrons trigger a Townsend avalanche, resulting in charge amplification by a factor on the order of 10^3 , Sect. 2.4.

Due to the large electron mobility, the gas amplification process is finished in less than 1 ns. Charge signals start to evolve on the anode $q_a(t)$ and the mesh $q_m(t)$ as soon as electrons reach the anode, the temporal behavior of the major signal component is caused by the drift of positive ions towards the mesh. The observable signals are given by the sum of charge carriers from gas amplification on the electrode and the induction signal by charge carriers in the amplification region

$$\begin{aligned} q_a(t) &= q_e(t) - q_{ai}(t) \\ q_m(t) &= q_i(t) - q_{mi}(t) . \end{aligned} \tag{2.26}$$

$q_e(t)$ is the constant charge of electrons, produced in gas amplification, $q_{ai}(t)$ is the negative mirror charge of the positive ion distribution in the amplification gap on the anode, $q_i(t)$ is the positive ion charge on the mesh and $q_{mi}(t)$ is the negative mirror charge on the mesh, produced by positive ions in the amplification region.

The four components and the resulting detectable charge signals are shown in Fig. 2.7. Details of the underlying calculation can be found in Bortfeldt (2010).

Unlike a low-energy X-ray, a charged particles creates a trace of ionization charge in the drift gap, the electrons enter the amplification region over a period on the order of 100 ns, depending on their drift velocity, diffusion and the drift gap width. Each electron creates a signal with the described temporal behavior. The resulting

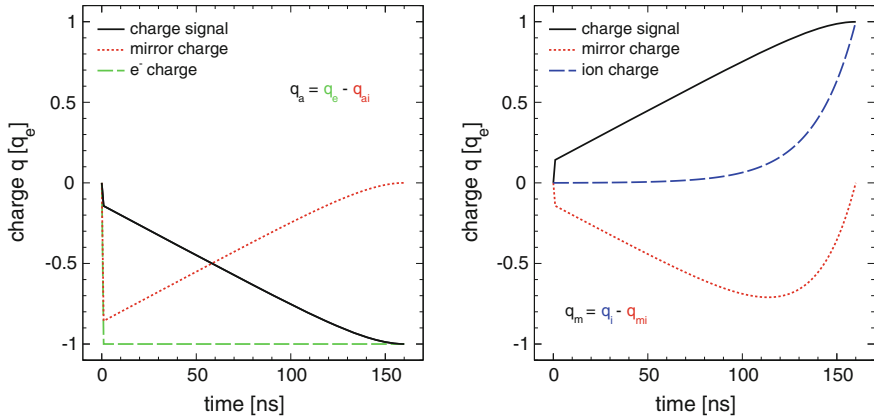


Fig. 2.7 Calculated charge signals on anode *left* and mesh *right*, assuming an amplification gap of 128 μm , an amplification field of 43.0 kV/cm and a detector gas mixture of Ar:CO₂ 93:7 vol% (Bortfeldt et al. 2012)

observable charge signal is determined by the ionization charge distribution, the electron drift velocity and the single electron signal evolution time, that is given by the maximum drift time of positive ions from the anode to the mesh. In order to decrease the overall signal length, the electron as well as the ion drift time has to be reduced.

The electron drift time can be controlled by a variation of the drift field, up to a certain extent. Also the use of modified Ar:CO₂ gas mixtures allows for a variation, see Fig. 2.4. Considerable variations however can only be achieved by using different gas admixtures such as methane or isobutane.

The ion drift velocity can only feebly be adjusted by the amplification field, since the desired gas gain defines the amplification field. The ion drift is furthermore rather insensitive to gas admixtures. In order to considerably increase the ion drift velocity, completely different gas constituents have to be used. Following Table 2.2, the use of lighter noble gases such as neon or helium allows for a considerable reduction of the signal length.

2.6 Mesh Transparency

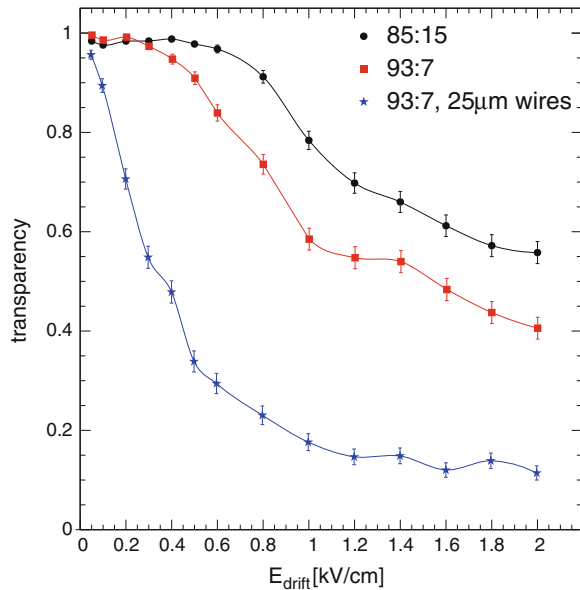
Electrons from ionization processes in the drift region move in the homogeneous field between cathode and mesh towards the mesh. Upon entering the amplification region, avalanches are created, which can be detected by the readout electronics. The signal rise time is defined by the sum of the electron drift time in the drift region and the ion drift time in the amplification region. The electron drift time $t_d = d_d/v_d$, can reach values above several 100 ns at low fields, depending on the

drift gap width d_d and the drift velocity v_d , see Fig. 2.4. If the integration time of the applied preamplifier electronics is smaller than the electron drift time, only a fraction of the ionization charge is detected, leading to smaller pulse heights at small drift fields. Furthermore, recombination with positive ions and formation of negative ions hinders the electrons from reaching the amplification region at low fields. Especially contaminations with oxygen lead to a strong attachment, due to its high electronegativity. Strictly speaking, these low-field effects are not correlated to the electron mesh transparency, although all effects, leading to the observed drift field dependence of the pulse height are often subsumed in the term mesh transparency, that is used as synonym for the transparency of the mesh for electrons.

The optical transparency of a micro-mesh with 400 lines per inch (lpi) and $18\text{ }\mu\text{m}$ wire diameter is 0.513, and only 0.368 for a wire diameter of $25\text{ }\mu\text{m}$. Fortunately, the transparency of the mesh for electrons is much higher, as the high electric amplification field extends into the mesh holes and guides electrons around the wires into the high field region. The electron mesh transparency is strongly correlated with the transverse electron diffusion, see Fig. 2.5, since diffusion off the electric field lines leads to a loss of electrons on the mesh wires. Furthermore, the ratio of the drift and the amplification field defines the extension of the high field into the low field region. With increasing drift field, more and more field lines end on the mesh, leading to a reduction of detectable charge.

In Fig. 2.8, the electron mesh transparency for two different Ar:CO₂ gas mixtures and two different mesh configurations is shown. Drift of single electrons with subsequent gas amplification has been simulated with the GARFIELD++ (Veenhof 2010)

Fig. 2.8 Electron mesh transparency as a function of the drift field, calculated with GARFIELD for a woven stainless steel mesh with $63.5\text{ }\mu\text{m}$ pitch and $18\text{ }\mu\text{m}$ wire diameter for two different Ar:CO₂ gas mixtures at $20\text{ }^\circ\text{C}$ and 1013 mbar and for a woven mesh with equal pitch, but $25\text{ }\mu\text{m}$ wire diameter. An amplification field of $E_{\text{amp}} = 39.1\text{ kV/cm}$ is assumed. The simulation is described in detail in Moll (2013), results courtesy of P. Lösel



package. The simulation has been developed and validated for the investigation of discharges in Micromegas, induced by strongly ionizing particles (Moll 2013).

For low drift fields, a transparency of 0.95–0.99 is reached. Clearly visible is the influence of the increasing drift field on the electron mesh transparency. The discussed and usually observed initial increase at very low fields is not correctly simulated, as contamination through oxygen and the integration time of the preamplifier electronics are not taken into account. The thin-wire mesh has a transparency plateau close to 1 for drift fields up to $E_{\text{drift}} \sim 0.3 \text{ kV/cm}$ for the 93:7 vol% and up to $E_{\text{drift}} \sim 0.6 \text{ kV/cm}$ for the 85:15 vol% gas mixture. The difference between the two gas mixtures is caused by the different transverse electron diffusion, cf. Fig. 2.5. The observed decrease of the transparency for the thin-wire mesh is sometimes less pronounced than in the simulation, since the wire mesh is often rolled or pressed to improve electron transparency.

For the thick-wire mesh, the loss of electrons with increasing drift field is much more pronounced, as the optical transparency already is considerably smaller.

Correlated to the electron transparency is the ion mesh transparency. It is often expressed by the so called ion backflow, which denotes the fraction of ions, produced in gas amplification, that are not neutralized at the mesh, but move into the drift region between mesh and cathode. As ion backflow limits the high-rate capabilities of Micromegas with very large drift regions such as Time-Projection-Chambers or at ultra-high rates, it is usually reduced as much as possible. In Sect. 5.4, a measurement of the ion backflow ratio in a resistive strip Micromegas is presented.

2.7 Discharges

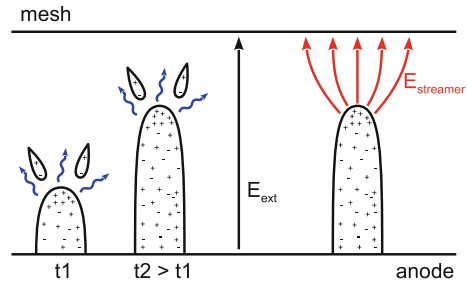
In Micromegas detectors non-destructive discharges between the mesh and the read-out structure are observed. A discharge leads to a conductive plasma between the mesh and the affected part of the anode structure and causes dead time, due to the necessary recharge of anode or mesh.

A discussion of discharge development measurements in different micro-pattern gaseous detectors can be found in Bressan et al. (1999).

The probability for spontaneous discharges, caused by dust or detector imperfections, can be strongly reduced by detector assembly under clean room conditions. Furthermore in non-resistive strip Micromegas a conditioning of the detector is possible, where dust, etc. is burned in air, prior to assembly, using an elevated high-voltage between mesh and anode.

Strongly ionizing particles such as alpha particles, nuclear fragments or low-energy protons from (n,p)-reactions can create a charge density in the amplification region sufficiently large, that a mesh directed streamer develops. A critical charge density of $1.77 \times 10^6 \text{ e}/0.01 \text{ mm}^2$ has been determined by Moll (2013) in a standard Micromegas with $128 \mu\text{m}$ amplification gap, operated with Ar:CO₂ 93:7 vol%.

Fig. 2.9 Streamer development in the amplification region of a Micromegas detector, following Raizer (1991)



The ionization induced streamer development in a planar electrode geometry has been discussed in detail by Raether (1964) and Raizer (1991). In the following a short, qualitative summary of the mechanism is given.

Under normal operation conditions, ionization charge, created by traversing particles, drifts into the high-field region between mesh and anode, where it is amplified. Positive ions from gas amplification drift within about 150 ns to the mesh, where they are neutralized. The total charge, created by a minimum ionizing particle and subsequent amplification is typically on the order of $6 \times 10^4 e$, too small to significantly alter the amplification field.

Strongly ionizing particles can create a much larger ion density in the amplification gap, drift of ions to the mesh competes with continuous ion production by ionization charge triggered gas avalanches. If the ion density in the proximity of the anode exceeds a critical value, the development of a mesh directed streamer begins: As the electric field of the positive ion cloud and its mirror charge on the anode is of the same order as the external amplification field, the resulting field between the upper part of the ion cloud, the so called streamer head, and the mesh is enhanced, between the ion cloud and the anode, the field is compensated. Photons from gas atom de-excitation and ion recombination leave the ion cloud and can ionize gas atoms in the vicinity. The photo-electrons trigger additional Townsend avalanches, that are directed towards the streamer head, due to the large electric field in the proximity of the streamer head, see $t = t_1$ in Fig. 2.9.

The ions produced in these secondary avalanches lead to a growth of the streamer towards the mesh, $t = t_2 > t_1$. Due to the compensation of the external electric field between the streamer head and the anode, a quasineutral plasma with rather low conductivity is present. It leads to an equalization of potentials between the anode and the streamer. As the streamer grows towards the mesh, the electric field between mesh and streamer head increases. At a certain point secondary electron emission from the mesh sets in, either due to ion-impact ionization or photo-effect. A larger number of electrons is created, that are accelerated towards the streamer and create a well conducting quasineutral plasma, that is on mesh potential as it touches the mesh. The large electric field in front of the head of this “back-streamer” leads to considerable gas amplification, such that the back-streamer develops towards the anode, leaving behind a well conducting plasma. Once the back streamer touches the anode the mesh discharges through the plasma channel onto the anode.

This discussion shows, that discharges in Micromegas, unlike in Geiger-Müller-tubes, are localized. Their impact on the detector performance can be greatly reduced by limiting the affected anode region. This can be achieved e.g. by covering the copper anode structure with a resistive layer (Sects. 3.1.2 and 3.1.3) or by coupling copper anode strips or pixels individually to high-voltage (Sect. 3.2).

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