Diplomacy Thesis

Simulation Studies at the MICROMEGAS Detector and Micropattern Applications in Medicine

Collaboration with C.E.R.N., Demokritos and N.T.U.A.

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Abstract

The intensive research aiming at matching the needs of creating more improved gaseous detectors helped in the emergence of more than a dozen of new ones. About in the beginning of the 1990 decade MICROMEGAS detector had been developed by Y. Giomataris and G. Charpak. MICROMEGAS is a gaseous detector based on amplification of electron avalanches in short gaps of the order of 100 μm at atmospheric pressure. It has specific properties of stability with respect to the gap length or gas pressure connected with the region of electric field at which it operates. It combines a good energy resolution, 5.4% FWHM at 22 keV X-rays, excellent position resolution that can reach 12 μm; high rate capability well adapted to particle physics experiments. A simple pre-amplification opens the way for orders of magnitude improvement in rate, and, in some cases, in accuracy in beta ray detection. The ease of construction is illustrated by the availability, for some experiments, of large detectors. The properties of MICROMEGAS open the way for promising applications in radiology, X-ray imaging with intense sources, imaging of neutrons, beta- and gamma-rays.
Περίληψη

Η διεξοδική έρευνα προκειμένου να δημιουργηθούν ανιχνευτές αερίου πιο βελτιωμένοι βοήθησε στη δημιουργία χιλιάδων νέων ανιχνευτών. Περίπου στις αρχές της δεκαετίας του '90 αναπτύχθηκε από τον Ι. Γιοματάρη και G. Charpak ο MICROMEGAS ανιχνευτής. Ο MICROMEGAS είναι λοιπόν ένας ανιχνευτής αερίου, του οποίου η λειτουργία βασίζεται στην ενίσχυση της χιονοστιβάδας των ηλεκτρονίων σε μικρό κενό της τάξεως των 100 μμ σε ατμοσφαιρική πίεση. Έχει συγκεκριμένες ιδιότητες ως προς την ομοιομορφία όσον αφορά το μήκος του κενού ή της πίεσης του αερίου που συνδέεται με την περιοχή του ηλεκτρικού πεδίου στο οποίο και λειτουργεί. Επιπλέον έχει καλή ενεργειακή διακριτική ικανότητα, 5.4% FWHM σε ακτίνες X των 22 keV, υψηλό χωρική διακριτική ικανότητα η οποία μπορεί να φτάσει τα 12 μμ- ικανότητα υψηλού ρυθμού μπορεί να αποκτηθεί για πειράματα σωματιδιακής φυσικής. Μια απλή προενίσχυση μπορεί να βελτιώσει κατά ένα μέγεθος το ρυθμό και, σε μερικές περιπτώσεις, την ακρίβεια ανίχνευσης της β-ακτινοβολίας. Η απλή δομή της δίνει την ευκολία στη δημιουργία ανιχνευτών μεγάλων διαστάσεων και αυτό βοηθάει σε μερικά πειράματα καθώς επίσης στη βιομηχανία και στην ιατρική. Οι ιδιότητες του MICROMEGAS ανιχνευτή ανοίγει το δρόμο για πολλές εφαρμογές του στην ακτινολογία, την ανίχνευση ακτινών X σκληρής δέσμης, καθώς και την ανίχνευση νετρονίων, β- και γ- ακτινοβολιών.
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Introduction

The intensive research aiming at matching the needs of high luminosity colliders has helped in the emergence of more than a dozen of new detectors. Some have similar structures, some rely on different approaches. While the choice of the detectors for the internal trackers for LHC (see Appendix A) is a closed matter, with the selection of the semiconductor detectors, chosen mostly out of reasons of prudence and not cost, the diversified research has resulted in new instruments which can be of decisive importance in many fields: physics, astrophysics, biology, medicine [1].

Multiwire proportional chambers have been originally designed for high-rate applications [2]. Their flux capability was mainly limited by the positive-ion space charge created because of the low ion drift velocity with a typical drift time of several tenths of microseconds. Their spatial resolution was limited by the wire spacing, which was of the order of 1 mm [3].

To overcome these limitations, a new technique, the microstrip gas chamber (MSGC), has been developed over the last eighteen years. Wires are replaced by strips printed on an insulating support; a high electric field region, sufficient for electron multiplication, is created between the thin cathode and anode conductive strips. It is new class of gas detector relying on the microelectronics technology. The small inter-strip pitch allows a good spatial resolution, inferior to 100 μm and the fast collection of the charges offers the possibility to cope with higher counting rates. One limitation of the MSGC detector is the fact that the avalanche multiplication does not exceed $10^4$, because of breakdown on the
insulator surface. Positive ions created during the avalanche process and accumulated on
the insulator, locally modify the electric field and cause a drop of the gain in the
irradiated area of the detector [4]. A lot of effort has been invested, during the last ten
years, resolve the charging-up problems by a careful chaise of the resistivity of the
substrate or a special treatment of its surface. Another type in this class of gas detectors,
the micro-gap chamber [5], was recently developed aiming to resolve the charging-up
problem and giving superior results in terms of rate capability and spatial resolution [3].

Another possible way out is the use of a special asymmetric configuration of the
multiwire structure [6, 7] with alternating anodes and field-shaping wires, mounted close
to the cathode plane with engraved pick-up strips orthogonal to the wire direction. The
performance of this structure equals that of the MSGC in terms of rate capability and
spatial resolution. In addition it can achieve higher electron multiplication factors and it
operates in a stable fashion for long irradiation periods. The drawback here is the use of
delicate wires and the wire stretching force, which is proportional to the total number of
wires acting in the wire frame; this therefore has to be of substantial thickness [3].

We are going to present a new approach where the wire plane is replaced by a thin
electroformed micromesh. The amplification occurs between the mesh plane and the
microstrip plane. A small gap, of about 100 \( \mu \text{m} \), between the anode and cathode plane is
kept by precise insulating spacers. The device operates as a two-stage parallel plate
avalanche chamber and it is called MICROMEGAS (MICRO-MEsh-GAseous Structure)
[3].

The design of MICROMEGAS is offering, for several applications, substantial
advantages in energy, space and time resolution, microscopic granularity on large
surfaces, insensitivity to discharges, simplicity of construction and capacity to identify
and reduce some sources of background, which can be of great interest in the search for
rare events [1].
Some of the properties of this structure and their advantageous consequences in a series of experiments related to charged particle localization, neutron imaging, gamma or X-ray imaging will be described [1]. For example, it is already used in CERN Experiments. CAST Experiment (see Appendix B) has three X-ray detectors which one of them is MICROMEGAS. Moreover, this detector is going to be used in ATLAS Experiment upgrade in order a high performance large TCP to be built, instead of using MWPC (Multi Wire Proportional Champsers) new Micro-Pattern Detector (MPGD) readout [8].
Basic Nuclear Processes in Radioactive Sources

Radioactive sources [1] provide a convenient means of testing and calibrating detectors and are essential tools. An understanding of the basic nuclear processes in radioactive sources is a necessity. We shall begin by briefly reviewing these processes and describing the characteristics of the resulting radiations.

Nuclei can undergo a variety of processes resulting in the emission of radiation of some form. We can divide the processes into two categories: radioactivity and nuclear reactions. The radiation emitted in both of these processes may be electromagnetic or corpuscular. The electromagnetic radiations consist of x-rays and y-rays while the corpuscular emissions include α-particles, β-electrons and positrons, internal conversion electrons, Auger electrons, neutrons, protons, and fission fragments, among others.

Table 1.1 summarizes some of the more common types of radiation found in laboratory sources. Each radiation type is characterized by an energy spectrum which is indicative of the nuclear process underlying it.
1.1. Nuclear Level Diagrams

The use of nuclear energy level diagrams is very important and we are going to provide a compact and convenient way of representing the changes which occur in nuclear transformations. These are usually plotted in the following way. For a given nucleus with atomic number $Z$ and mass $A$, the energy levels are plotted as horizontal lines on some arbitrary vertical scale. The spin and parity of each of these states may also be indicated. Keeping the same mass number $A$, the energy levels of neighbouring nuclei ($Z - 1, A$), ($Z + 1, A$), ... are now also plotted on this energy scale with $Z$ ordered in the horizontal direction, as illustrated in Fig. 1.1.

<table>
<thead>
<tr>
<th>Type</th>
<th>Origin</th>
<th>Process</th>
<th>Charge</th>
<th>Mass [MeV]</th>
<th>Spectrum (energy)</th>
</tr>
</thead>
<tbody>
<tr>
<td>a-particles</td>
<td>Nucleus</td>
<td>Nuclear decay</td>
<td>+ 2</td>
<td>3727.33</td>
<td>Discrete [MeV]</td>
</tr>
<tr>
<td>$\beta^-$-rays</td>
<td>Nucleus</td>
<td>Nuclear decay</td>
<td>-1</td>
<td>0.511</td>
<td>Continuous [keV-MeV]</td>
</tr>
<tr>
<td>$\beta^+$-rays (positrons)</td>
<td>Nuclear</td>
<td>Nuclear decay</td>
<td>+ 1</td>
<td>0.511</td>
<td>Continuous [keV-MeV]</td>
</tr>
<tr>
<td>$\gamma$-rays</td>
<td>Nucleus</td>
<td>Nuclear deexcitation</td>
<td>0</td>
<td>0</td>
<td>Discrete [keV-MeV]</td>
</tr>
<tr>
<td>x-rays</td>
<td>Electron cloud</td>
<td>Atomic deexcitation</td>
<td>0</td>
<td>0</td>
<td>Discrete [eV-keV]</td>
</tr>
<tr>
<td>Internal conversion electrons</td>
<td>Electron cloud</td>
<td>Nuclear deexcitation</td>
<td>-1</td>
<td>0.511</td>
<td>Discrete [high keV]</td>
</tr>
<tr>
<td>Auger electrons</td>
<td>Electron cloud</td>
<td>Atomic deexcitation</td>
<td>-1</td>
<td>0.511</td>
<td>Discrete [eV-keV]</td>
</tr>
<tr>
<td>Neutrons</td>
<td>Nucleus</td>
<td>Nuclear reaction</td>
<td>0</td>
<td>939.57</td>
<td>Continuous or discrete [keV-MeV]</td>
</tr>
<tr>
<td>Fission fragments</td>
<td>Nucleus</td>
<td>Fission</td>
<td>= 20</td>
<td>80-160</td>
<td>Continuous [30-150 MeV]</td>
</tr>
</tbody>
</table>
Figure 1.1. Nuclear level diagrams conveniently represent the transitions which can occur between nuclei.

For a system of $A$ nucleons, energy is represented on the vertical scale while atomic number is on the horizontal scale.

This reflects the fact that nuclei with different $Z$ but the same $A$ may simply be treated as different states of a system of $A$ nucleons. The relation of the energy levels of a nucleus $(Z, A)$ to other nuclei in the same $A$ system is therefore made apparent.

With the exception of $\alpha$-decay, radioactive decay may now be viewed as simply a transition from a higher energy state to a lower energy state within the same system of $A$ nucleons. For example, consider the $\beta$-decay process, which will be discussed in the next section. This reaction involves the decay

$$ (Z, A) \rightarrow (Z + 1, A) + e^- + \bar{v} $$

where the final state in the nucleus $(Z + 1, A)$ may be the ground state or some excited state. This is shown in Fig. 1.1 by the arrow descending to the right. The atomic number $Z$ increases by one, but $A$ remains constant. The changes that occur can be immediately seen: for example, the energy available for the reaction as given by the difference in height between the two levels, the spin and parity changes, etc. If it is possible for the same initial state to make transitions to several different final states, this can also be represented by several arrows emanating from the initial state to the various possible final
states. The relative probability of each decay branch (i.e., the branching ratio) may also be indicated next to the corresponding arrow.

In a similar way, transitions which follow the first may also be diagrammed. For example, suppose the final state of the above example is an excited state of the \((Z+1, A)\) nucleus, it may then make a gamma transition to the ground state or to another excited state. This type of transition is indicated by a vertical line since \(Z\) remains unchanged \((\text{see Fig. 1.1})\). Other transitions, such as a further \(\beta\)-decay or some other process, may be represented in a similar manner. In this way, the types of radiation emitted by a particular radioactive source and their origins may be easily displayed.

### 1.2. Alpha Decay

Alpha particles are \(^4\)He nuclei, i.e., a bound system of two protons and two neutrons, and are generally emitted by very heavy nuclei containing too many nucleons to remain stable. The emission of such a nucleon cluster as a whole rather than the emission of single nucleons is energetically more advantageous because of the particularly high binding energy of the \(\alpha\)-particle. The parent nucleus \((Z,A)\) in the reaction is thus transformed via

\[
(Z, A) \rightarrow (Z - 2, A - 4) + \alpha
\]

Theoretically, the process was first explained by Gamow and Condon and by Gurney as the tunneling of the \(\alpha\)-particle through the potential barrier of the nucleus. Alpha particles, therefore, show a monoenergetic energy spectrum. As well, since barrier transmission is dependent on energy, all \(\alpha\)-sources are generally limited to the range \(4 - 6 MeV\) with the higher energy sources having the higher transmission probability and thus the shorter half-life. For this reason also, most \(\alpha\)-decays are directly to the ground state of the daughter nucleus since this involves the highest energy change. Decays to excited states of the daughter nucleus are nevertheless possible, and in such nuclei, the energy spectrum shows several monoenergetic lines each corresponding to \(\alpha\)-
decay to one of these states. Some of the more commonly used sources are listed below in Table 1.2.

Because of its double charge, $+2e$, alpha particles have a very high rate of energy loss in matter. The range of a $5 \text{ MeV}$ $\alpha$-particle in air is only a few centimeters, for example. For this reason it is necessary to make $\alpha$-sources as thin as possible in order to minimize energy loss and particle absorption. Most $\alpha$-sources are made, in fact, by depositing the isotope on the surface of a suitable backing material and protecting it with an extremely thin layer of metal foil.

**Table 1.2. Characteristics of some alpha emitters**

<table>
<thead>
<tr>
<th>Isotope</th>
<th>Half-life</th>
<th>Energies [MeV]</th>
<th>Branching</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{241}\text{Am}$</td>
<td>433 y</td>
<td>5.486</td>
<td>85%</td>
</tr>
<tr>
<td></td>
<td></td>
<td>5.443</td>
<td>12.8%</td>
</tr>
<tr>
<td>$^{241}\text{Po}$</td>
<td>138 d</td>
<td>5.305</td>
<td>100%</td>
</tr>
<tr>
<td>$^{241}\text{Cm}$</td>
<td>163 d</td>
<td>6.113</td>
<td>74%</td>
</tr>
<tr>
<td></td>
<td></td>
<td>6.070</td>
<td>26%</td>
</tr>
</tbody>
</table>

**1.3. Beta Decay**

Beta particles are fast electrons or positrons which result from the weak-interaction decay of a neutron or proton in nuclei which contain an excess of the respective nucleon. In a neutron-rich nucleus, for example, a neutron can transform itself into a proton via the process

$$n \rightarrow p + e^- + \bar{\nu}$$

where an electron and anti neutrino are emitted. (The proton remains bound to the nucleus.) The daughter nucleus now contains one extra proton so that its atomic number is increased by 1. Similarly, in nuclei with too many protons, the decay

$$p \rightarrow n + e^+ + \nu$$
can occur, where a positron and a neutrino are now emitted and the atomic number is decreased by 1. Both are mediated by the same weak interaction.

A basic characteristic of the $\beta$-decay process is the continuous energy spectrum of the $\beta$-particle. This is because the available energy for the decay (the $Q$-value) is shared between the $\beta$-particle and the neutrino (or antineutrino) which usually goes undetected. An energy spectrum from $^{201}$Bi is shown in Fig. 1.2. If the small recoil energy of the daughter nucleus is ignored, the maximum energy of this spectrum should correspond to the $Q$-value for the reaction. For most beta sources, this maximum value ranges from a few tens of keV to a few MeV.

In very many $\beta$-sources, the daughter nucleus is left in an excited state which decays immediately with the emission of one or more $\gamma$ photons (see §1.5). This is illustrated in the level diagram shown in Fig. 1.3. These sources, therefore, are also emitters of $\gamma$ radiation. Most $\beta$-sources are of this type. Pure $\beta$-emitters exist but the list is astonishingly short as is seen in Table 1.3.

![Energy spectrum of beta decay electrons from $^{210}$Bi](Figure 1.2)

![Nuclear level diagrams of a few common gamma sources](Figure 1.3)
Some $\beta$-sources may also have more than one decay branch, i.e., they can decay to different excited states of the daughter nucleus. Each branch constitutes a separate $\beta$-decay with an end-point energy corresponding to the energy difference between the initial and final states and is in competition with the other branches. The total $\beta$-spectrum from such a source is then a superposition of all the branches weighted by their respective decay probabilities.

Since electrons lose their energy relatively easily in matter, it is important that $\beta$-sources be thin in order to allow the $\beta$ to escape with a minimum of energy loss and absorption. This is particularly important for positron sources since the positron can annihilate with the electrons in the source material or surrounding container. A too thick $p^+$ source will exhibit a distorted $p^+$ spectrum and an enormous background of 511 keV annihilation photons.

**Table 1.3.** List of pure $\beta^-$ emitters

<table>
<thead>
<tr>
<th>Source</th>
<th>Half-life</th>
<th>$E_{\text{max}}$ [MeV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^3$H</td>
<td>12.26 y</td>
<td>0.0186</td>
</tr>
<tr>
<td>$^{14}$C</td>
<td>5730 y</td>
<td>0.156</td>
</tr>
<tr>
<td>$^{32}$P</td>
<td>14.28 d</td>
<td>1.710</td>
</tr>
<tr>
<td>$^{33}$P</td>
<td>24.4 d</td>
<td>0.248</td>
</tr>
<tr>
<td>$^{35}$S</td>
<td>87.9 d</td>
<td>0.167</td>
</tr>
<tr>
<td>$^{36}$Cl</td>
<td>$3.08 \times 10^5$ y</td>
<td>0.714</td>
</tr>
<tr>
<td>$^{45}$Ca</td>
<td>165 d</td>
<td>0.252</td>
</tr>
<tr>
<td>$^{63}$Ni</td>
<td>92 y</td>
<td>0.067</td>
</tr>
<tr>
<td>$^{90}$Sr$/^{90}$Y</td>
<td>27.7 y/ 64 h</td>
<td>0.546/ 2.27</td>
</tr>
<tr>
<td>$^{99}$Tc</td>
<td>$2.12 \times 10^5$ y</td>
<td>0.292</td>
</tr>
<tr>
<td>$^{147}$Pm</td>
<td>2.62 y</td>
<td>0.224</td>
</tr>
<tr>
<td>$^{204}$Tl</td>
<td>3.81 y</td>
<td>0.766</td>
</tr>
</tbody>
</table>
1.4. Electron Capture (EC)

As an alternative to $p^+$ emission, proton-rich nuclei may also transform themselves via the capture of an electron from one of the atomic orbitals

$$e^- + p \rightarrow n + \nu.$$  \hspace{1cm} (1.5)

This reaction is essentially the same as $p^+$-decay but with the $p$ particle transposed to the left side. The nuclear level diagram for EC is therefore identical to that for $p^+$ emission. Since only the neutrino is emitted, electron capture would seem to be a reaction almost impossible to observe, given the well-known difficulty of detecting such a particle! The capture of the electron, however, leaves a hole in the atomic shell which is filled by another atomic electron, giving rise to the emission of a characteristic x-ray or Auger electrons (see Appendix C). These radiations are, of course, much more amenable to detection and can be used to signal the capture reaction. In general, it is the $K$ electron which is most likely captured, although, $L$-capture is also possible but with a much smaller probability.

1.5. Gamma Emission

Like the electron shell structure of the atom, the nucleus is also characterized by discrete energy levels. Transitions between these levels can be made by the emission (or absorption) of electromagnetic radiation of the correct energy, i.e., with an energy equal to the energy difference between the levels participating in the transition. The energies of these photons, from a few hundred $keV$ to a few $MeV$, characterize the high binding energy of nuclei. These high-energy photons were historically named $\gamma$-rays, and, like atoms, show spectral lines characteristic of the emitting nucleus. Level diagrams illustrating the specific energy structure of some typical $\gamma$-ray sources are shown in Fig. 1.3.

Most $\gamma$-sources are “placed” in their excited states as the result of a $\beta$-disintegration, although excited nuclear states are often created in nuclear reactions also. Since electrons and positrons are more easily absorbed in matter, the $p\beta$-particles in such sources can be
“filtered” out by enveloping them with sufficient absorbing material, leaving only the more penetrating $\gamma$-ray.

### 1.1.1. Isomeric States

Although most excited states in nuclei make almost immediate transitions to a lower state, some nuclear states may live very much longer. Their de-excitation is usually hindered by a large spin difference between levels (i.e., a *forbidden* transition) resulting in lifetimes ranging from seconds to years. A nuclide which is “trapped” in one of these metastable states will thus show radioactive properties different from those in more normal states. Such nuclei are called *isomers* and are denoted by an $m$ next to the mass number in their formulae, e.g. $^{60m}\text{Co}$ or $^{69m}\text{Zn}$.
This chapter concerns the basic reactions [1] which occur when radiation encounters matter and the effects produced by these processes. As will be seen in the following chapters, these processes are the basis of all current particle detection devices and thus determine the sensitivity and efficiency of a detector.

2.1. Preliminary Notions and Definitions

To open our discussion of radiation in matter, we first review a few basic notions concerning the interaction of particles.

2.1.1. Cross Section

The collision or interaction of two particles is generally described in terms of the *cross section*. This quantity essentially gives a measure of the probability for a reaction to occur and may be calculated if the form of the basic interaction between the particles is known. Formally, the cross-section is defined in the following manner. Consider a beam of particles $I$ incident upon a target particle 2 as shown in Fig. 2.1. Assume that the beam is much broader than the target and that the particles in the beam are uniformly distributed in space and time. We can then speak of a *flux* of $F$ incident particles per unit
area per unit time. Now look at the number of particles scattered\(^1\) into the solid angle \(d\Omega\) per unit time. Because of the randomness of the impact parameters, this number will fluctuate over different finite periods of measuring time. However, if we average many finite measuring periods, this number will tend towards a fixed \(d\sigma / d\Omega\), where \(N_s\) is the average number scattered per unit time. The **differential cross section** is then defined as the ratio

\[
\frac{d\sigma}{d\Omega}(E, \Omega) = \frac{1}{F} \frac{dN_s}{d\Omega},
\]

that is, \(d\sigma / d\Omega\) is the average fraction of the particles scattered into \(d\Omega\) per unit time per unit flux \(F\). In terms of a single quantum mechanical particle, this may be reformulated as the scattered probability current in the angle \(d\Omega\) divided by the total incident probability passing through a unit area in front of the target.

Note that because of the dimensions of \(P\), \(d\sigma\) has dimensions of area, which leads to the heuristic interpretation of \(d\sigma\) as the geometric cross sectional area of the target intercepting the beam. That fraction of the flux incident on this area will then obviously interact while all those missing \(d\sigma\) will not. This is only a visual aid, however, and should in no way be taken as a real measure of the physical dimensions of the target.

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{scattering_diagram.png}
\caption{Definition of the scattering cross section}
\end{figure}

---

\(^1\) By *scattering* here, we mean *any* reaction in which an outgoing particle is emitted into \(\Omega\). The incident particle need not retain its identity.
In general, the value of $d\Omega$ will vary with the energy of the reaction and the angle at which the particle is scattered. We can calculate a total cross section for any scattering whatsoever at an energy $E$ defined as the integral of $d\sigma / d\Omega$ over all solid angles,

$$\sigma(E) = \int d\Omega \frac{d\sigma}{d\Omega}.$$  \hspace{1cm} (2.2)

While the above example is easily visualized, it is not a practical case. In real situations, of course, the target is usually a slab of material containing many scattering centers and it is desired to know how many interactions occur on the average. Assuming that the target centers are uniformly distributed and the slab is not too thick so that the likelihood of one center sitting in front of another is low, the number of centers per unit perpendicular area which will be seen by the beam is then $N \delta x$ where $N$ is the density of centers and $\delta x$ is the thickness of the material along the direction of the beam. If the beam is broader than the target and $A$ is the total perpendicular area of the target, the number of incident particles which are eligible for an interaction is then $F A$. The average number scattered into $d\Omega$ per unit time is then

$$N_s(\Omega) = FA\delta x \frac{d\sigma}{d\Omega}.$$ \hspace{1cm} (2.3)

The total number scattered into all angles is similarly

$$N_{tot} = FA N \delta x \sigma.$$ \hspace{1cm} (2.4)

If the beam is smaller than the target, then we need only set $A$ equal to the area covered by the beam. Then $FA \rightarrow n_{inc}$, the total number of incident particles per unit time. In both cases, now, if we divide (2.4) by $FA$, we have the probability for the scattering of a single particle in a thickness $\delta x$,

$$\text{Prob. of interaction in } \delta x = N\sigma \delta x.$$ \hspace{1cm} (2.5)

This is an important quantity and we will come back to this probability later.

### 2.1.2. Interaction Probability in a Distance $x$. Mean Free Path

In the previous section, we discussed the probability for the interaction of a particle traveling through a thin slab of matter containing many interaction centers. Let us consider the more general case of any thickness $x$. To do this, we ask the opposite question: what is the probability for a particle not to suffer an interaction in a distance $x$?
This is known as the \textit{survival} probability and may be calculated in the following way. Let \( P(x) \): probability of \textit{not} having an interaction after a distance \( x \), \( w \, dx \): probability of having an interaction between \( x \) and \( x + dx \). The probability of \textit{not} having an interaction between \( x \) and \( x + dx \) is then
\[
P(x + dx) = P(x)(1 - w \, dx),
\]
\[
P(x) + \frac{dP}{dx} \, dx = P - P \, w \, dx,
\]
\[
dP = -w \, P \, dx,
\]
\[
P = C \exp(-wx),
\]
where \( C \) is a constant. Requiring that \( P(0) = 1 \), we find \( C = 1 \). The probability of the particle surviving a distance \( x \) is thus exponential in distance. From this, of course, we see immediately that the probability of suffering an interaction \textit{anywhere} in the distance \( x \) is just
\[
P_{\text{surv}}(x) = 1 - \exp(-wx),
\]
while the probability of the particle suffering a collision between \( x \) and \( x + dx \) after surviving the distance \( x \) is
\[
F(x) \, dx = \exp(-wx)w \, dx.
\] (2.8)

Now let us calculate the mean distance, \( \lambda \), travelled by the particle without suffering a collision. This is known as the \textit{mean free path}. Thus,
\[
\lambda = \frac{\int x \, P(x) \, dx}{\int P(x) \, dx} = \frac{1}{w}.
\] (2.9)

Intuitively, \( \lambda \) must be related to the density of interaction centers and the cross-section, for as we have seen, this governs the probability of interaction. To find this relation, let us return to our slab of material. For a small thickness \( \delta x \), the interaction probability (2.7) can then be approximated as
\[
P_{\text{surv}} = 1 - \left(1 - \frac{\delta x}{\lambda} + \ldots\right) = \frac{\delta x}{\lambda}
\] (2.10)
where we have expanded the exponential and kept only the first order term. Comparing with (2.5), we find,
so that our survival probability becomes

\[ P(x) = \exp \left( \frac{-x}{\lambda} \right) = \exp(-N\sigma x) \]  

and the interaction probabilities

\[ P_{\text{int}}(x) = 1 - \exp \left( \frac{-x}{\lambda} \right) = 1 - \exp(-N\sigma x), \]

\[ F(x) = \exp \left( \frac{-x}{\lambda} \right) \frac{dx}{\lambda} = \exp(-N\sigma x) N\sigma dx. \]

### 2.1.3. Surface Density Units

A unit very often used for expressing thicknesses of absorbers is the *surface density* or *mass thickness*. This is given by the mass density of the material times its thickness in normal units of length, i.e.,

\[ \text{mass thickness} = \rho \cdot t \]

with \( \rho \): mass density, \( t \): thickness, which, of course, yields dimensions of mass per area, e.g. g/cm\(^2\).

For discussing the interaction of radiation in matter, mass thickness units are more convenient than normal length units because they are more closely related to the density of interaction centers. They thus have the effect of normalizing materials of differing mass densities. As will be seen later, equal mass thicknesses of different materials will have roughly the same effect on the same radiation.

### 2.2. Energy Loss of Heavy Charged Particles by Atomic Collisions

In general, two principal features characterize the passage of charged particles through matter: (1) a loss of energy by the particle and (2) a deflection of the particle from its incident direction. These effects are primarily the result of two processes:

1. inelastic collisions with the atomic electrons of the material
2. elastic scattering from nuclei.
These reactions occur many times per unit path length in matter and it is their cumulative result which accounts for the two principal effects observed. These, however, are by no means the only reactions which can occur. Other processes include

3. emission of Cherenkov radiation
4. nuclear reactions
5. bremsstrahlung.

In comparison to the atomic collision processes, they are extremely rare, however, and with the exception of Cherenkov radiation, will be ignored in this treatment.

For reasons which will become clearer in the following sections, it is necessary to separate charged particles into two classes: (1) electrons and positrons, and (2) heavy particles, i.e., particles heavier than the electron. This latter group includes the muons, pions, protons, a-particles and other light nuclei. Particles heavier than this, i.e., the heavy ions, although technically part of this latter group, are excluded in this discussion because of additional effects which arise.

Of the two electromagnetic processes, the inelastic collisions are almost solely responsible for the energy loss of heavy particles in matter. In these collisions ($\sigma = 10^{-17} - 10^{-16}$ cm$^2$), energy is transferred from the particle to the atom causing an ionization or excitation of the latter. The amount transferred in each collision is generally a very small fraction of the particle’s total kinetic energy; however, in normally dense matter, the number of collisions per unit path length is so large, that a substantial cumulative energy loss is observed even in relatively thin layers of material. A 10 MeV proton, for example, already loses all of its energy in only 0.25 mm of copper! These atomic collisions are customarily divided into two groups: soft collisions in which only an excitation results, and hard collisions in which the energy transferred is sufficient to cause ionization. In some of the hard reactions, enough energy is, in fact, transferred such that the electron itself causes substantial secondary ionization. These high-energy recoil electrons are sometimes referred to as $\delta$-rays or knock-on electrons.
Elastic scattering from nuclei also occurs frequently but not as often as electron collisions. In general very little energy is transferred in these collisions since the masses of the nuclei of most materials are usually large compared to the incident particle. In cases where this is not true, for example, an a-particle in hydrogen, some energy is also lost through this mechanism. Nevertheless, the major part of the energy loss is still due to atomic electron collisions.

The inelastic collisions are, of course, statistical in nature, occurring with a certain quantum mechanical probability. However, because their number per macroscopic pathlength is generally large, the fluctuations in the total energy loss are small and one can meaningfully work with the average energy loss per unit path length. This quantity, often called the stopping power or simply \( dE/dx \), was first calculated by Bohr using classical arguments and later by Bethe-Bloch and others using quantum mechanics [2].

2.3. Energy Loss of Electrons and Positrons

Like heavy charged particles, electrons and positrons also suffer a collisional energy loss when passing through matter. However, because of their small mass an additional energy loss mechanism comes into play: the emission of electromagnetic radiation arising from scattering in the electric field of a nucleus (bremsstrahlung). Classically, this may be understood as radiation arising from the acceleration of the electron (or positron) as it is deviated from its straight-line course by the electrical attraction of the nucleus. At energies of a few \( MeV \) or less, this process is still a relatively small factor. However, as the energy is increased, the probability of bremsstrahlung quickly shoots up so that at a few 10’s of \( MeV \), loss of energy by radiation is comparable to or greater than the collision-ionization loss. At energies above this critical energy, bremsstrahlung dominates completely.

The total energy loss of electrons and positrons, therefore, is composed of two parts:

\[
\left( \frac{dE}{dx} \right)_{\text{tot}} = \left( \frac{dE}{dx} \right)_{\text{rad}} + \left( \frac{dE}{dx} \right)_{\text{coil}}.
\]  

(2.16)
2.4. Energy Straggling: The Energy Loss Distribution

Our discussion of energy loss up until now has been concerned mainly with the *mean* energy loss suffered by charged particles when passing through a thickness of matter. For any given particle, however, the amount of energy lost will *not*, in general, be equal to this mean value because of the statistical fluctuations which occur in the number of collisions suffered and in the energy transferred in each collision. An initially monoenergetic beam, after passing through a fixed thickness of material, will therefore show a distribution in energy rather than a delta-function peak shifted down by the mean energy loss as given by the $dE/dx$ formula. We have already seen these fluctuations in the form of range straggling. This, in fact, is the same problem viewed from a different angle: instead of observing the fluctuations in energy loss for a fixed thickness of absorber, we observe the fluctuations in thickness of pathlength for a fixed loss in energy.

From a theoretical point of view, calculating the distribution of energy losses for a given thickness of absorber is a difficult mathematical problem and is generally divided into two cases: thick absorbers and thin absorbers.

2.5. The Interaction of Photons

The behavior of photons in matter (in our case, x-rays and $\gamma$-rays) is dramatically different from that of charged particles. In particular, the photon’s lack of an electric charge makes impossible the many inelastic collisions with atomic electrons so characteristic of charged particles. Instead, the main interactions of x-rays and $\gamma$-rays in matter are:

1. Photoelectric Effect
2. Compton Scattering (including Thomson and Rayleigh Scattering)
3. Pair Production.

Also possible, but much less common, are nuclear dissociation reactions, for example, $(\gamma, n)$, which we will neglect in our discussion.
These reactions explain the two principal qualitative features of x-rays and y-rays: (1) x-rays and \(\gamma\)-rays are many times more penetrating in matter than charged particles, and (2) a beam of photon’s is \textit{not} degraded in energy as it passes through a thickness of matter, only attenuated in intensity. The first feature is, of course, due to the much smaller cross section of the three processes relative to the inelastic electron collision cross section. The second characteristic, however, is due to the fact the three processes above remove the photon from the beam entirely, either by absorption or scattering. The photons which pass straight through, therefore, are those which have not suffered any interactions at all. They therefore retain their original energy. The total number of photons is, however, reduced by the number which has interacted. The attenuation suffered by a photon beam can be shown, in fact, to be exponential with respect to the thickness, i.e.,

\[
I(x) = I_0 \exp(-\mu x),
\]

with \(I_0\): incident beam intensity; \(x\): thickness of absorber; \(\mu\): absorption coefficient.

The absorption coefficient is a quantity which is characteristic of the absorbing material and is directly related to the total interaction cross-section. This is a quantity often referred to when discussing y-ray detectors. However, let us first discuss the three processes individually before turning to the calculation of the absorption coefficient.

\subsection*{2.5.1. Photoelectric Effect}

The photoelectric effect involves the absorption of a photon by an atomic electron with the subsequent ejection of the electron from the atom. The energy of the outgoing electron is then

\[
E = h\nu - B.E.
\]

where \(B.E.\) is the binding energy of the electron.

Since a free electron cannot absorb a photon and also conserve momentum, the photoelectric effect always occurs on bound electrons with the nucleus absorbing the recoil momentum. Figure 2.2 shows a typical photoelectric cross section as a function of incident photon energy. As can be seen, at energies above the highest electron binding
energy of the atom (the \( K \) shell), the cross section is relatively small but increases rapidly as the \( K \)-shell energy is approached. Just after this point, the cross section drops drastically since the \( K \)-electrons are no longer available for the photoelectric effect. This drop is known as the \( K \) absorption edge. Below this energy, the cross section rises once again and dips as the \( L \), \( M \), levels, etc. are passed. These are known respectively as the \( L \)-absorption edges, \( M \)-absorption edge, etc.

![Figure 2.2. Calculated photoelectric cross section for lead](image)

Theoretically, the photoelectric effect is difficult to treat rigorously because of the complexity of the Dirac wavefunctions for the atomic electrons. For photon energies above the \( K \)-shell, however, it is almost always the \( K \) electrons which are involved. If this is assumed and the energy is nonrelativistic, i.e., \( h\nu \ll m_e c^2 \), the cross-section can then be calculated using a Born approximation. In such a case, one obtains

\[
\Phi_{\text{photon}} = 4 a^4 \sqrt{2} Z^5 \phi_0 \left( m_e c^2 / h\nu \right)^{7/2} \text{ per atom}
\]

where \( \phi_0 = 8\pi r_e^2 / 3 = 6.651 \times 10^{-25} \text{ cm}^2 \); \( a = 1/137 \).

For energies closer to the \( K \)-edge, (2.19) must be multiplied by a correction factor to give
\[
\Phi_{\text{photon}} = \frac{2^7 \pi (137)^3}{Z^2} \left[ \frac{v_k^4}{\nu} \right] \exp \left( -4\xi \cot^{-1}\xi \right) \frac{1 - \exp(-2\pi \xi)}{\nu}, \text{ per atom} \quad (2.20)
\]

Where \( h\nu_k = (Z - 0.03) \frac{1}{2} m_e c^2 a^2 / 2 \) and \( \xi = \sqrt{v_k / \left( v - v_k \right)} \). For \( \nu \) very close to \( v_k \), \( \xi \gg 1 \), so that (2.20) can be simplified to

\[
\Phi_{\text{photon}} = \frac{6.3 \times 10^{-18}}{Z^2} \left( \frac{v_k}{\nu} \right)^{8/3}. \quad (2.21)
\]

Formulas for the \( L \) and \( M \) shells have also been calculated, but these are even more complicated than those above [3].

It is interesting to note the dependence of the cross section on the atomic number \( Z \). This varies somewhat depending on the energy of the photon, however, at \( MeV \) energies, this dependence goes as \( Z \) to the 4th or 5th power. Clearly, then, the higher \( Z \) materials are the most favored for photoelectric absorption, and, as will be seen in later chapters, are an important consideration when choosing \( \gamma \)-ray detectors.

2.5.2. Compton Scattering

Compton scattering is probably one of the best understood processes in photon interactions. As will be recalled, this is the scattering of photons on free electrons. In matter, of course, the electrons are bound; however, if the photon energy is high with respect to the binding energy, this latter energy can be ignored and the electrons can be considered as essentially free.

Figure 2.3 illustrates this scattering process. Applying energy and momentum conservation, the following relations can be obtained.

\[
h\nu' = \frac{h\nu}{1 + \gamma \left( 1 - \cos \theta \right)}.
\]
Figure 2.3. Kinematics of Compton scattering

\[ T = h \nu - h \nu' = h \nu \cdot \frac{\gamma(1 - \cos \theta)}{1 + \gamma(1 - \cos \theta)}, \]

\[ \cos \theta = 1 - \frac{2}{(1 + \gamma)^2 \tan^2 \phi + 1}, \] (2.22)

\[ \cos \theta = (1 + \gamma) \tan \frac{\theta}{2}, \]

where \( \gamma = h \nu / m_e c^2 \). Other relations between the various variables may be found by substitution in the above formulae.

The cross section for Compton scattering was one of the first to be calculated using quantum electrodynamics and is known as the Klein-Nishina formula:

\[ \frac{d\sigma}{d\Omega} = \frac{r_e^2}{2} \left[ \frac{1}{1 + \gamma(1 - \cos \theta)} \right]^2 \left[ 1 + \cos^2 \theta + \frac{\gamma^2(1 - \cos^2 \theta)^2}{1 + \gamma(1 - \cos^2 \theta)} \right] \] (2.23)

where \( r_e \) is the classical electron radius. Integration of this formula over \( d\Omega \), then, gives the total probability per electron for a Compton scattering to occur,

\[ \sigma_e = 2 \pi r_e^2 \left[ \frac{1 + \gamma}{\gamma^2} \left( \frac{2(1 + \gamma)}{1 + 2\gamma} - \frac{1}{\gamma} \ln(1 + 2\gamma) \right) + \frac{1}{2\gamma} \ln(1 + 2\gamma) - \frac{1 + 3\gamma}{(1 + 2\gamma)^2} \right]. \] (2.24)

Figure 2.4 plot this total cross section as a function of energy.
Two useful quantities which can be calculated from the Klein-Nishina formula are the Compton scattered and Compton absorption cross sections. The Compton scattered cross section, $\sigma^s$, is defined as the average fraction of the total energy contained in the scattered photon, while the absorption cross section, $\sigma^a$, is the average energy transferred to the recoil electron. Since the electron is stopped by the material, this is the average energy-fraction absorbed by the material in Compton scattering. Obviously, the sum must be equal to $\sigma_c$

$$\sigma_c = \sigma^s + \sigma^a$$  \hspace{1cm} (2.25)

To calculate $\sigma^s$, we form

$$\frac{d\sigma^s}{d\Omega} = \frac{h \nu' d\sigma}{h \nu} d\Omega,$$  \hspace{1cm} (2.26)

which after integration yields

$$\sigma^s = \pi r_e^2 \left[ \frac{1}{\gamma^3} \ln (1 + 2\gamma) + \frac{2(1 + \gamma)(2\gamma^2 - 2\gamma - 1)}{\gamma^2 (1 + 2\gamma)^2} + \frac{8\gamma^2}{3(1 + 2\gamma)^3} \right].$$  \hspace{1cm} (2.27)

![Figure 2.4. Total Compton scattering cross section](image-url)
The absorption cross section can then be simply calculated by
\[ \sigma^s = \sigma_c - \sigma^s. \] (2.28)

Another formula which we will make use of very often when discussing detectors is the energy distribution of the Compton recoil electrons. By substituting into the Klein-Nishina formula, one obtains
\[ \frac{d\sigma}{dT} = \frac{2 \pi r_e^2}{m_e c^2 \gamma^2} \left[ 2 + \frac{s^2}{\gamma^2 (1-s)^2} + \frac{s}{1-s} \left( s - \frac{2}{\gamma} \right) \right], \] (2.29)
where \( s = \frac{T}{h \nu} \). Figure 2.5 shows this distribution for several incident photon energies. The maximum recoil energy allowed by kinematics is given by
\[ T_{\text{max}} = h \nu \left( \frac{2 \gamma}{1 + 2 \gamma} \right) \] (2.30)
(see Eq. 2.22) and is known as the Compton edge.

![Figure 2.5. Energy distribution of Compton recoil electrons. The sharp drop at the maximum recoil energy is known as the Compton edge](image)
2.5.2.1. Thomson and Rayleigh Scattering

Related to Compton scattering are the classical processes of Thomson and Rayleigh scattering. Thomson scattering is the scattering of photons by free electrons in the classical limit. At low energies with respect to the electron mass, the Klein-Nishina formula, in fact, reduces to the Thomson cross-section,

\[
\sigma = \frac{8\pi}{3} r_e^2.
\]  

(2.31)

Rayleigh scattering, on the other hand, is the scattering of photons by atoms as a whole. In this process, all the electrons in the atom participate in a coherent manner. For this reason it is also called coherent scattering.

In both processes, the scattering is characterized by the fact that no energy is transferred to the medium. The atoms are neither excited nor ionized and only the direction of the photon is changed. At the relatively high energies of x-rays and γ-rays, Thomson and Rayleigh scattering are very small and for most purposes can be neglected.

2.5.3. Pair Production

The process of pair production involves the transformation of a photon into an electron-positron pair. In order to conserve momentum, this can only occur in the presence of a third body, usually a nucleus. Moreover, to create the pair, the photon must have least an energy of 1.022 MeV.

Theoretically, pair production is related to bremsstrahlung by a simple substitution rule, so that once the calculations for one process are made, results for the other immediately follow. As for bremsstrahlung, the screening by the atomic electrons surround, g the nucleus plays an important role in pair production. The cross sections are thus dependent on the parameter \(\xi\) [see Appendix D –Eq. (D.38)], which is now defined by

\[
\xi = \frac{100 m_e c^2 \hbar v}{E_+ E_- Z^{1/3}}
\]  

(2.32)

with \(E_+\): total energy of outgoing positron; \(E_-\): total energy of outgoing electron.
At extreme relativistic energies and arbitrary screening, a Born approximation calculation gives the formula

\[
d\tau = 4Z^2r^2a \frac{dE}{(h\nu)} \left\{ \left( E_+^2 + E_0^2 \right) \left[ \frac{\phi_1(\xi)}{4} - \frac{1}{3} \ln Z - f(Z) \right] + \frac{2}{3} E_+E_0 \left[ \frac{\phi_2(\xi)}{4} - \frac{1}{3} \ln Z - f(Z) \right] \right\}
\]

(2.33)

where \( \phi_1 \) and \( \phi_2 \) are the screening functions used in (D.40) and the other variables are as defined in (D.39).

As before, this formula simplifies in the limiting cases of no screening and complete screening. Thus for no screening \( (\xi \gg 1) \), we obtain

\[
d\tau = 4Z^2a_r^2 \frac{dE}{(h\nu)} \left\{ E_+^2 + E_0^2 + \frac{(2E_+ E_0)}{3} \left[ \ln \left( \frac{2E_+ E_0}{h\nu m_e c^2} \right) - \frac{1}{2} - f(Z) \right] \right\},
\]

(2.34)

while for complete screening, \( \xi \to 0 \),

\[
d\tau = 4Z^2a_r^2 \frac{dE}{(h\nu)} \left\{ \left( E_+^2 + E_0^2 \right) \left[ \ln \left( \frac{183Z^{-1/3}}{f(Z)} \right) - \frac{E_+ E_0}{9} \right] \right\}.
\]

(2.35)

Because of the Born approximation, these formulae are not very accurate for high \( Z \) or low energy\(^2\).

To obtain the total pair production cross section, a numerical integration of the above expressions must generally be performed. In the case of no screening with \( m_e c^2 \ll h\nu \ll 137m_e c^2 Z^{-1/3} \), an analytic integration is possible yielding

\[
\tau_{\text{pair}} = 4Z^2a_r^2 \left[ \frac{7}{9} \left( \ln \frac{2h\nu}{m_e c^2} - f(Z) \right) - \frac{109}{54} \right].
\]

(2.36)

\(^2\) A more complicated formula valid for low energies and no screening has been derived by Bethe and Heitler and is given in the article by Bethe and Ashkin [4] along with a somewhat simpler formula from Hough.
Similarly for complete screening, $h\nu \gg 137 m_e c^2 Z^{-1/3}$,

$$
\tau_{\text{pair}} = 4 Z^2 a^2 c^2 \left[ \frac{7}{9} \left( \ln \left( 183 Z^{-1/3} \right) - f(Z) \right) - 1/54 \right].
$$

(2.37)

For all other cases, a numerical integration of (2.33) must be performed. Figure 2.6 illustrates the energy dependence of the total pair cross section.

As for bremsstrahlung, pair production may also occur in the field of an atomic electron. Not surprisingly, a similar result is obtained for the cross section, but smaller by about a factor $Z$. To approximately account for this interaction, then, one need only replace $Z^2$ by $Z(Z+)$ in the above formulae.

From the total cross section, it is interesting to calculate the mean free path, $\lambda_{\text{pair}}$ of a $\gamma$-ray for pair production. Thus, using (2.37)

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure2_6.png}
\caption{Experimental values for the photon total-absorption cross section as a function of energy and of $Z$. The solid lines are calculated values, with the pair-production contribution given by the Bethe-Heitler theory corrected for Born-approximation failure in heavy elements.}
\end{figure}
where \( N \) is the density of atoms and we have ignored the small constant term. This may be recognized as being very similar to the radiation length, and, in fact, comparison with (D.52) shows
\[
\lambda_{\text{pair}} = \frac{9}{7} L_{\text{rad}}. \tag{2.39}
\]

### 2.5.4. The Total Absorption Coefficient and Photon Attenuation

The total probability for a photon interaction in matter is the sum of the individual cross sections outlined above. If we calculate the cross-section per atom, this yields
\[
\sigma = \Phi_{\text{photon}} + Z \sigma_c + \tau_{\text{pair}}, \tag{2.40}
\]
where we have multiplied the Compton cross-section by \( Z \) to take into account the \( Z \) electrons per atom. This is shown in Fig. 2.7 for the ease of lead. If we now multiply \( \sigma \) by the density of atoms, \( N \), we then obtain the probability per unit length for an interaction,
\[
\mu = N \sigma = \sigma \left( N_a \frac{\rho}{A} \right) \tag{2.41}
\]
with \( N_a \): Avogadro’s Number; \( \rho \): density of the material; \( A \): molecular weight.

This is more commonly known as the total absorption coefficient and is just the inverse of the mean free path of the photon. From (2.12), then, it follows that the fraction of photons surviving a distance \( x \) is then
\[
I/I_0 = \exp(-\mu x), \tag{2.42}
\]
where \( I_0 \) is the incident intensity.

For compounds and mixtures, the total absorption coefficient may be calculated using Bragg’s rule (D.23),
\[
\frac{\mu}{\rho} = w_1 \frac{\mu_1}{\rho_1} + w_2 \frac{\mu_2}{\rho_2} + \ldots, \tag{2.43}
\]
where \( w_i \) is the weight fraction of each element in the compound.
Like the photon, the neutron lacks an electric charge, so that it is not subject to Coulomb interactions with the electrons and nuclei in matter. Instead, its principal means of interaction is through the strong force with nuclei. These reactions are, of course, much rarer in comparison because of the short range of this force. Neutrons must come within \( \approx 10^{-13} \, cm \) of the nucleus before anything can happen, and since normal matter is mainly empty space, it is not surprising that the neutron is observed to be a very penetrating particle.

When the neutron does interact, however, it may undergo a variety of nuclear processes depending on its energy. Among these are:
1. Elastic scattering from nuclei, i.e., \( A(n, n)A \). This is the principal mechanism of energy loss for neutrons in the MeV region.

2. Inelastic scattering, e.g., \( A(n, n')A^* \), \( A(n, n)A \), etc. In this reaction, the nucleus is left in an excited state which may later decay by gamma-ray or some other form of radiative emission. In order for the inelastic reaction to occur, the neutron must, of course, have sufficient energy to excite the nucleus, usually on the order of 1 MeV or more. Below this energy threshold, only elastic scattering may occur.

3. Radiative neutron capture, i.e., \( n + (Z, A) \rightarrow \gamma + (Z, A+1) \). In general, the cross-section for neutron capture goes approximately as \( 1/u \) where \( u \) is the velocity of neutron. Absorption is most likely, therefore, at low energies. Depending on element, there may also be resonance peaks superimposed upon this \( 1/u \) dependence. At these energies, of course, the probability of neutron capture is greatly enhanced.

4. Other nuclear reactions, such as \((n, p)\), \((n, d)\), \((n, a)\), \((n, t)\), \((n, a p)\), etc. in which neutron is captured and charged particles are emitted. These generally occur in eV to keV region. Like the radiative capture reaction, the cross section generally falls as \( 1/u \). Resonances may also occur depending on the element.

5. Fission, i.e., \((n, f)\). Again this is most likely at thermal energies.

6. High energy hadron shower production. This occurs only for very high energy neutrons with \( E > 100 MeV \).

Because of the strong energy dependence of neutron interactions, it has become customary to classify neutrons according to their energy, although no specific boundaries are prescribed between classes. In general, high energy neutrons are considered to be those with energies above \( \approx 100 MeV \) or so, whereas those between a few ten’s of MeV and a few hundred keV are known as fast neutrons. Between \( \approx 100 keV \) and \( \approx 0.1 eV \), where nuclear resonance reactions occur, neutrons are referred to epithermal. At lower
energies comparable to the thermal agitation energy at room temperature, (i.e., \( E = kT = 1/40 \text{eV} \)), neutrons are known as thermal or slow. Going even lower to energies of million or micro-\( \text{eV} \), neutrons come under the appellation cold or ultra-cold.

The total probability for a neutron to interact in matter is given by the sum of individual cross sections, i.e.,

\[
\sigma_{\text{tot}} = \sigma_{\text{elastic}} + \sigma_{\text{inelastic}} + \sigma_{\text{capture}} + \ldots.
\]  

(2.44)

Figure 2.8 gives an example of the total reaction cross-section for neutrons on a few materials versus neutron energy. Here the energy dependence is quite smooth.

If we multiply (2.44) by the density of atoms we can obtain the mean free path length

\[
\frac{1}{\lambda} = N \sigma_{\text{tot}} = \frac{N_a \rho}{A} \sigma_{\text{tot}}.
\]

(2.45)

**Figure 2.8.** Total reaction cross-sections for neutrons in water, paraffin and photons [5]
In analogy to photons, then, a beam of neutrons passing through matter will be exponentially attenuated

\[ N = N_0 \exp\left(-\frac{x}{\lambda}\right) \]  

(2.46)

where \( x \) is the thickness of the material. Equation (2.46), of course, is useful only for a collimated beam of neutrons. For the more general case of a noncollimated source, a sophisticated transport equation is usually necessary.
Chapter 3

Ionization Detectors

As an introduction to the following chapter, we will define and describe here some general characteristics common to detectors. Moreover a short description of the phenomenology used in gaseous detectors is introduced.

3.1. General Characteristics of Detectors

3.1.1. Sensitivity

The first consideration for a detector [1] is its sensitivity, i.e., its capability of producing a usable signal for a given type of radiation and energy. No detector can be sensitive to all types of radiation at all energies. Instead, they are designed to be sensitive to certain types of radiation in a given energy range. Going outside this region usually results in an unusable signal or greatly decreased efficiency.

Detector sensitivity to a given type of radiation of a given energy depends on several factors:

1. the cross section for ionizing reactions in the detector
2. the detector mass
3. the inherent detector noise
4. the protective material surrounding the sensitive volume of the detector.

### 3.1.2. Detector Response

In addition to detecting the presence of radiation, most detectors are also capable of providing some information on the energy of the radiation. This follows since the amount of ionization produced by radiation in a detector is proportional to the energy it loses in the sensitive volume. If the detector is sufficiently large such that the radiation is completely absorbed, then this ionization gives a measure of the energy of the radiation. Depending on the design of the detector, this information may or may not be preserved as the signal is processed, however [1].

#### 3.1.3. Energy Resolution; Fano Factor

For detectors which are designed to measure the energy of the incident radiation, the most important factor is the energy resolution [1]. This is the extent to which the detector can distinguish two close lying energies. In general, the resolution can be measured by sending a monoenergetic beam of radiation into the detector and observing the resulting spectrum. Ideally, of course, one would like to see a sharp delta-function peak. In reality, this is never the case and one observes a peak structure with a finite width, usually Gaussian in shape. This width arises because of fluctuations in the number of ionizations and excitations produced.

The resolution is usually given in terms of the full width at half maximum of the peak (FWHM). Energies which are closer than this interval are usually considered unresolvable; this is illustrated in Fig. 3.1. If we denote this width as \( \Delta E \), then the relative resolution at the energy \( E \) is

\[
\text{Resolution} = \frac{\Delta E}{E}
\]

Eq. (3.1) is usually expressed in percent. A NaI detector has about a 8% or 9% resolution for \( \gamma \)-rays of about 1 MeV, for example, while germanium detectors have resolutions on the order of 0.1%.
Figure 3.1. Definition of energy resolution. Two peaks are generally considered to be resolved if they are separated by a distance greater than their full widths at half maximum (FWHM). The solid line shows the sum of two identical Gaussian peaks separated by just this amount.

In general, the resolution is a function of the energy deposited in the detector, with the ratio (3.1) improving with higher energy. This is due to the Poisson or Poisson-like statistics of ionization and excitation. Indeed, it is found that the average energy required to produce an ionization is a fixed number, \( w \), dependent only on the material. For a deposited energy \( E \), therefore, one would expect on the average, \( J = E / w \) ionizations. Thus as energy increases, the number of the ionization events also increases resulting in smaller relative fluctuations.

To calculate the fluctuations it is necessary to consider two cases. For a detector in which the radiation energy is not totally absorbed, for example, a thin transmission detector which just measures the \( dE/dx \) loss of a passing particle, the number of signal producing reactions is given by a Poisson distribution. The variance is then given by

\[
\sigma^2 = J ,
\]

where \( J \) is the mean number of events produced. The energy dependence of the resolution can then be seen to be

\[
R = 2.35 \frac{\sqrt{J}}{J} = 2.35 \frac{\sqrt{w}}{E} ,
\]

(3.3)
where the factor 2.35 relates the standard deviation of a Gaussian to its FWHM. Thus the resolution varies inversely as the square root of the energy.

If the full energy of the radiation is absorbed as is the case for detectors used in spectroscopy experiments, the naive assumption of Poisson statistics is incorrect. And indeed, it is observed that the resolution of many such detectors is actually smaller than that calculated from Poisson statistics. The difference here is that the total energy deposited is a fixed, constant value, while in the previous case, the energy deposited can fluctuate. The total number of ionizations which can occur and the energy lost in each ionization is thus constrained by this value. Statistically, this means that the ionization events are not all independent so that Poisson statistics is not applicable. Fano [2] was the first to calculate the variance under this condition and found

$$\sigma^2 = F J$$  \hspace{1cm} (3.4)

where $J$ is the mean ionization produced and $F$ is a number known as the Fano factor. The factor $F$ is a function of all the various fundamental processes which can lead to an energy transfer in the detector. This includes all reactions which do not lead to ionization as well, for example, phonon excitations, etc. It is thus an intrinsic constant of the detecting medium. Theoretically, $F$ is very difficult to calculate accurately as it requires a detailed knowledge of all the reactions which can take place in the detector. From (3.4), the resolution is then given by

$$R = 2.35 \sqrt{\frac{F J}{J}} = 2.35 \sqrt{\frac{F w}{E}}.$$  \hspace{1cm} (3.5)

3.1.4. The Response Function

For the measurement of energy spectra, an important factor which must be considered is the response function of the detector for the type of radiation being detected. This is the spectrum of pulse heights observed from the detector when it is bombarded by a monoenergetic beam of the given radiation. The response function [1] of a detector at a
given energy is determined by the different interactions which the radiation can undergo in the detector and its design and geometry.

### 3.1.5. Response Time

A very important characteristic of a detector is its response time [1]. This is the time which the detector takes to form the signal after the arrival of the radiation. This is crucial to the timing properties of the detector. For good timing, it is necessary for the signal to be quickly formed into a sharp pulse with a rising flank as close to vertical as possible. In this way a more precise moment in time is *marked* by the signal.

The duration of the signal is also of importance. During this period, a second event cannot be accepted either because the detector is insensitive or because the second signal will *pile up* on the first. This contributes to the *dead* time of the detector and limits the count rate at which it can be operated.

### 3.1.6. Detector Efficiency

Two types of efficiency [1] are generally referred to when discussing radiation detection: *absolute* efficiency and *intrinsic* detection efficiency. The absolute or total efficiency of a detector is defined as that fraction of events emitted by the source which is actually registered by the detector, i.e.,

\[
\epsilon_{\text{tot}} = \frac{\text{events registered}}{\text{event emitted by source}}.
\]

(3.6)

This is a function of the detector geometry and the probability of an interaction in the detector.

### 3.2. Ionization and Transport Phenomena in Gas

Because of the importance of ionization detectors much work has been and still is devoted to the ionization process and the movement of electrons and ions in gases. We shall, therefore, devote some time to reviewing some of these processes in the following sections.
3.2.1. Ionizing Collisions

A charged particle that traverses the gas of a drift chamber leaves a track of ionization along its trajectory. The encounters with the gas atoms are purely random and are characterized by a mean free flight path $\lambda$ [3] between ionizing encounters given by the ionization cross-section per electron $\sigma_1$ and the density $N$ of electrons:

$$\lambda = 1/(N \sigma_1)$$  \hspace{1cm} (3.7)

Therefore, the number of encounters along any length $L$ has a mean of $L/A$, and the frequency distribution is the Poisson distribution

$$P(L/A, k) = \frac{(L/\lambda)^k}{k!} \exp(-L/\lambda)$$  \hspace{1cm} (3.8)

It follows that the probability distribution $f(l)dl$ of the free flight paths $l$ between encounters is an exponential, because the probability of finding zero encounters in the interval $l$ times the probability of one encounter in $dl$ is equal to

$$f(l)dl = P(l/\lambda, 0)P(dl/\lambda, 1) = (1/\lambda)\exp(-l/\lambda)dl.$$  \hspace{1cm} (3.9)

From (3.4) we obtain the probability of having zero encounters along a track length $L$:

$$P(L/\lambda, 0) = \exp(-L/\lambda).$$  \hspace{1cm} (3.10)

Equation (3.10) provides a method for measuring $\lambda$. If a gas counter with sensitive length $L$ is set up so that the presence of even a single electron in $L$ will always give a signal, then its inefficiency may be identified with expression (3.10), thus measuring $\lambda$.

In Table 3.1 we present additional measurements of a larger number of gases that are employed in drift chambers. These primary ionization cross-sections $\sigma_p$ were measured by Rieke and Prepejchal [4] in the vicinity of the minimum at different values of $\gamma$ and interpolated to the minimum $\sigma_{p}^{\text{min}}$ at $\gamma^{\text{min}}$, using the parametrization of the Bethe-Bloch formula (see Sect. D.2.1). The mean free path $\lambda$ is related to $\sigma_p$ by the number density $N_m$ of molecules:
\[ \lambda = 1 / \left( N_m \sigma_p \right). \quad (3.11) \]

The measurement errors are within \( \pm 4\% \) (see the original paper for details). In comparison with the values presented in Table 1.1, the measurements are in rough agreement, except for argon.

### 3.2.2. Different Ionization Mechanisms

We distinguish between primary and secondary ionization [3]. In primary ionization, one or sometimes two or three electrons are ejected from the atom \( A \) encountered by the fast particle, say a \( \pi \) meson:

\[ \pi A \rightarrow \pi A^+ e^- , \pi A^+ e^- e^- , \ldots \quad (3.12) \]

**Table 3.1.** Minimal primary ionization cross-sections \( \sigma_p \) for charged particles in some gases, and relativistic velocity factor \( \gamma_{\text{min}} \) of minimum, according to measurements done by Rieke and Prepejchal [4]

<table>
<thead>
<tr>
<th>Gas</th>
<th>( \sigma_p \left( 10^{-20} \text{ cm}^2 \right) )</th>
<th>( \gamma_{\text{min}} )</th>
<th>Gas</th>
<th>( \sigma_p \left( 10^{-20} \text{ cm}^2 \right) )</th>
<th>( \gamma_{\text{min}} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>H(_2)</td>
<td>18.7</td>
<td>3.81</td>
<td>i-C(<em>4)H(</em>{10})</td>
<td>333</td>
<td>3.56</td>
</tr>
<tr>
<td>He</td>
<td>18.6</td>
<td>3.68</td>
<td>n-C(<em>3)H(</em>{12})</td>
<td>434</td>
<td>3.56</td>
</tr>
<tr>
<td>Ne</td>
<td>43.3</td>
<td>3.39</td>
<td>neo-C(<em>5)H(</em>{12})</td>
<td>433</td>
<td>3.45</td>
</tr>
<tr>
<td>Ar</td>
<td>90.3</td>
<td>3.39</td>
<td>n-C(<em>6)H(</em>{14})</td>
<td>526</td>
<td>3.51</td>
</tr>
<tr>
<td>Xe</td>
<td>172</td>
<td>3.39</td>
<td>C(_2)H(_2)</td>
<td>126</td>
<td>3.60</td>
</tr>
<tr>
<td>O(_2)</td>
<td>92.1</td>
<td>3.43</td>
<td>C(_2)H(_4)</td>
<td>161</td>
<td>3.58</td>
</tr>
<tr>
<td>CO(_2)</td>
<td>132</td>
<td>3.51</td>
<td>CH(_3)OH</td>
<td>155</td>
<td>3.65</td>
</tr>
<tr>
<td>C(_2)H(_6)</td>
<td>161</td>
<td>3.58</td>
<td>C(_2)H(_2)OH</td>
<td>230</td>
<td>3.51</td>
</tr>
<tr>
<td>C(_3)H(_8)</td>
<td>269</td>
<td>3.47</td>
<td>(CH(_3))(_2)CO</td>
<td>277</td>
<td>3.54</td>
</tr>
</tbody>
</table>

Most of the charge along a track is from secondary ionization where the electrons are ejected from atoms not encountered by the fast particle. This happens either in collisions of ionization electrons with atoms,

\[ e^- A \rightarrow e^- A^+ e^- , e^- A^+ e^- e^- , \ldots \quad (3.13) \]
or through intermediate excited states $A^*$. An example is the following chain of reactions involving the collision of the excited state with a second species, $B$, of atoms or molecules that is present in the gas:

$$\pi A \rightarrow \pi A^*$$

or

$$e^- A \rightarrow e^- A^*, \quad (3.14)$$

$$A^* B \rightarrow AB^+ e^- \quad (3.15)$$

Reaction (3.15) occurs if the excitation energy of $A^*$ is above the ionization potential of B. In drift chambers, $A^*$ is often the metastable state of a noble gas created in reaction (3.14), and $B$ is one of the molecular additives (quenchers) that are required for the stability of proportional wire operation; $A^*$ may also be an optical excitation with a long lifetime due to resonance trapping. These effects are known under the names of Penning effect (involving metastables) and Jesse effect (involving optical excitations, also used more generally); obviously they depend very strongly on the gas composition and density.

Another example of secondary ionization through intermediate excitation has been observed in pure rare gases where an excited molecule $A^*_2$ has a stable ionized ground state $A^*_2^+$:

$$A^* A \rightarrow A^*_2 \rightarrow A^*_2^+ e^-. \quad (3.16)$$

The different contributions of processes (3.13-16) are in most cases unknown [5]. A pictorial summary of the processes discussed is given in Fig. 3.2.

Figure 3.3 shows the most probable number of electrons for 1 cm as a function of particle energy for several gas mixtures, otherwise e$^-$ ionization per cm for a muon. As we can see Ar mixtures have a higher stopping power than the Neon mixtures because Neon ($N = 10, A = 20$) is lighter than the Argon ($N = 18, A = 40$).
Figure 3.2. Pictorial classification of the ionization produced by a fast charged particle in a noble gas containing molecules with low ionization potential: (-) electron; (+) positive ion, single charge; (+ +) positive ion, double charge; (+) positive ion of the low-ionization species; (*) state excited above the lower ionization potential of the other species; ( ) (+) positive ion of noble gas molecule; – photon transmission, -- collision

Figure 3.3. Most probable number of electrons for 1 cm as a function of Particle energy for each gas mixture. Ionization e⁻/cm for μ (muon) by GARFIELD simulation program
3.2.3. Mean Number of Electron-Ion Pairs Created

Since the occurrence of the ionizing reactions above is statistical in nature, two identical particles will not, in general, produce the same number of ion-electron pairs. We can ask, however: What is the average number of ion-electron pairs (from all mechanisms) created for a given energy loss? Note that this is not equal to the energy loss divided by the ionization potential, since some energy is also lost to excitation! For gases, this average turns out to be on the order of 1 ion-electron pair per 30 eV of energy lost, that is, for a 3 keV particle, an average of 3000/30 = 100 ion-electron pairs will be created. Moreover, what is surprising is that this average value does not depend very strongly on particle type and only weakly on the type of gas. Table 3.2 gives a comparison of the measured values for this average for several types of gas used in ionization detectors.

Table 3.2. Excitation and ionization characteristics of various gases

<table>
<thead>
<tr>
<th>Gases</th>
<th>Excitation potential [eV]</th>
<th>Ionization potential [eV]</th>
<th>Mean energy for ion-electron pair creation [eV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>H₂</td>
<td>10.8</td>
<td>15.4</td>
<td>37</td>
</tr>
<tr>
<td>He</td>
<td>19.8</td>
<td>24.6</td>
<td>41</td>
</tr>
<tr>
<td>N₂</td>
<td>8.1</td>
<td>15.5</td>
<td>35</td>
</tr>
<tr>
<td>O₂</td>
<td>7.9</td>
<td>12.2</td>
<td>31</td>
</tr>
<tr>
<td>Ne</td>
<td>16.6</td>
<td>21.6</td>
<td>36</td>
</tr>
<tr>
<td>Ar</td>
<td>11.6</td>
<td>15.8</td>
<td>26</td>
</tr>
<tr>
<td>Kr</td>
<td>10.0</td>
<td>14.0</td>
<td>24</td>
</tr>
<tr>
<td>Xe</td>
<td>8.4</td>
<td>12.1</td>
<td>22</td>
</tr>
<tr>
<td>CO₂</td>
<td>10.0</td>
<td>13.7</td>
<td>33</td>
</tr>
<tr>
<td>CH₄</td>
<td></td>
<td>13.1</td>
<td>28</td>
</tr>
<tr>
<td>C₂H₂O</td>
<td></td>
<td>10.8</td>
<td>23</td>
</tr>
</tbody>
</table>

The average energy, \( w \), required for creating an electron-ion pair is important since it determines the efficiency and the energy resolution of the detector. From \( \left( \Delta E \right)^2 = \left( \Delta E_{\text{det}} \right)^2 + \left( \Delta E_{\text{elect}} \right)^2 + \ldots \), the resolution for a particle of energy \( E \) is
\[ R = 2.35 \sqrt{\frac{F w}{E}}. \]  

(3.17)

where \( F \) is the Fano factor for the gas medium, \( w \) the average energy required to produce an ionization is a fixed number. While the Fano factor is not well determined for most gases, it is clear that \( F \) is much less than 1. Table 3.3 gives some measured values for various noble gas mixtures.

**Table 3.3.** Excitation and ionization characteristics of various gases [6-8]

<table>
<thead>
<tr>
<th>Gas</th>
<th>( F )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ar 100%</td>
<td>0.2 ( \pm 0.01 )</td>
</tr>
<tr>
<td></td>
<td>0.2 ( \pm 0.02 )</td>
</tr>
<tr>
<td></td>
<td>&lt; 0.40 ( \pm 0.03 )</td>
</tr>
<tr>
<td>Ar + 80% Xe</td>
<td>&lt; 0.21 ( \pm 0.03 )</td>
</tr>
<tr>
<td>Ar + 24% Xe</td>
<td>&lt; 0.23 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Ar + 20% Xe</td>
<td>&lt; 0.16 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Ar + 5% Xe</td>
<td>&lt; 0.14 ( \pm 0.03 )</td>
</tr>
<tr>
<td>Ar + 5% Kr</td>
<td>&lt; 0.37 ( \pm 0.06 )</td>
</tr>
<tr>
<td>Ar + 20% Kr</td>
<td>&lt; 0.12 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Ar + 79% Kr</td>
<td>&lt; 0.13 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Xe 100%</td>
<td>&lt; 0.15 ( \pm 0.01 )</td>
</tr>
<tr>
<td></td>
<td>&lt; 0.15 ( \pm 0.03 )</td>
</tr>
<tr>
<td>Kr 100%</td>
<td>&lt; 0.23 ( \pm 0.01 )</td>
</tr>
<tr>
<td></td>
<td>&lt; 0.19 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Kr + 1.3% Xe</td>
<td>&lt; 0.19 ( \pm 0.01 )</td>
</tr>
<tr>
<td>Kr + 20% Xe</td>
<td>&lt; 0.21 ( \pm 0.02 )</td>
</tr>
<tr>
<td>Kr + 40% Xe</td>
<td>&lt; 0.22 ( \pm 0.01 )</td>
</tr>
<tr>
<td>Kr + 60% Xe</td>
<td>&lt; 0.21 ( \pm 0.01 )</td>
</tr>
<tr>
<td>Kr + 95% Xe</td>
<td>&lt; 0.21 ( \pm 0.01 )</td>
</tr>
</tbody>
</table>
3.3. Transport of Electrons and Ions in Gases

For ionization detectors, an understanding of the motion of the electrons and ions in gases is extremely important as these factors influence many operating characteristics of the detector. For the most part, this motion is described by the classical kinetic theory of gases. Two phenomena are of particular importance: diffusion, and drift in an electric field [1].

On the microscopic scale [2], the electrons or ions that drift through the gas are scattered on the gas molecules so that their direction of motion is randomized in each collision. On the average, they assume a constant drift velocity \( u \) in the direction of the electric field \( E \). The drift velocity \( u \) is much smaller than the instantaneous velocity \( c \) between collisions.

3.3.1. Diffusion

In the absence of an electric field, electrons and ions liberated by passing radiation diffuse uniformly outward from their point of creation. In the process they suffer multiple collisions with the gas molecules and lose their energy. They thus come quickly into thermal equilibrium with the gas and eventually recombine. At thermal energies, the velocities of the charges are described by the Maxwell distribution which gives a mean speed of

\[
\nu = \sqrt{\frac{8kT}{\pi m}},
\]

(3.18)

where \( k \) is Boltzmann’s constant, \( T \) the temperature and \( m \) the mass of the particle. Quite obviously, the average speed of the electrons is much greater than that of the ions due to their smaller mass. At room temperature, the electron speed is a few times \( 10^6 \) cm / s while the positive ion speeds are on the order of \( 10^4 \) cm / s.

From kinetic theory, the linear distribution of charges after diffusing a time \( t \) can be shown to be Gaussian,
\[
\frac{dN}{dx} = \frac{N_0}{\sqrt{4\pi Dt}} \exp\left(-\frac{x^2}{4Dt}\right),
\]

where \(N_0\) is the total number of charges, \(x\) the distance from the point of creation and \(D\) the diffusion coefficient. The rms spread in \(x\) is thus
\[
\sigma(x) = \sqrt{2Dt}.
\]

If three dimensions are considered, the spherical spread is given by
\[
\sigma(r) = \sqrt{6Dt},
\]
where \(r\) is the radial distance. The radial spread of ions in air under normal conditions, for example, is about 1 \(mm\) after 1 second [9]. The diffusion coefficient is a parameter which can be calculated from kinetic theory and can be shown to be
\[
D = \frac{1}{3} v \lambda
\]
where \(\lambda\) is the mean free path of the electron or ion in the gas. For a classical ideal gas, the mean free path is related to the temperature \(T\), and the pressure \(p\), by
\[
\lambda = \frac{1}{\sqrt{2}} \frac{kT}{\sigma_0 p},
\]
where \(\sigma_0\) is the total cross section for a collision with a gas molecule. Substituting (6.12) and (3.23) into (3.22) then gives the explicit expression
\[
D = \frac{2}{3\sqrt{\pi}} \frac{1}{p \sigma_0} \sqrt{\frac{(kT)^3}{m}}.
\]

The dependence of \(D\) on the various parameters of the gas now becomes evident.

Fig. 3.4 shows transverse diffusion coefficient as a function of the electric field for several gas mixtures, while Fig. 3.5 shows the longitudinal coefficient. What we observe from the Fig. 3.4 is that the CO\(_2\) is an efficient cooling agent in Ar mixtures. The ArCO\(_2\)CF\(_4\) mixture has low transverse diffusion and Transverse diffusion spreads the ionization electrons and can therefore help in making the signal broader of a MICROMEGAS. In Fig. 3.5, the longitudinal diffusion depends little on the composition for the electric fields found in the ionization region, no impact in MICROMEGAS.
Figure 3.4. Tranverse Diffusion as a function of electric field for several gas mixtures by GARFIELD simulation program

Figure 3.5. Longitudinal Diffusion as a function of Electric Field for several gas mixtures by GARFIELD simulation program
3.3.2. Drift of Electrons and Ions

In the presence of an electric field \([1]\), the electrons and ions freed by radiation are accelerated along the field lines towards the anode and cathode respectively. This acceleration is interrupted by collisions with the gas molecules which limit the maximum average velocity which can be attained by the charge along the field direction. The average velocity attained is known as the *drift velocity* of the charge and is superimposed upon its normal random movement. Compared to their thermal velocities, the drift speed of the ions is slow; however, for electrons this can be much higher since they are much lighter.

In kinetic theory, it is useful to define the *mobility* of a charge as

\[
\mu = \frac{u}{E}
\]

where \(u\) is the drift velocity and \(E\) the electric field strength. For positive ions, the drift velocity (Fig. 3.6) is found to depend linearly on the ratio \(E / p\), (also known as the reduced electric field), up to relatively high electric fields. At a constant pressure, this implies that the mobility \(\mu\) is a constant. For a given \(E\), it is also quite clear that \(\mu\) varies as the inverse of the pressure \(p\).

For ideal gases, in which the moving charges remain in thermal equilibrium, the mobility can be shown to be related to the diffusion constant by

\[
D / \mu = kT / e
\]

---

\(^3\) Let us consider an electron between two collisions. Because of its light mass, \(m\), the electron scatters isotropically and, immediately after the collision, it has forgotten any preferential direction. Some short time later, in addition to its instantaneous and randomly oriented velocity \(c\), the electron has picked up the extra velocity \(u\) equal to its acceleration along the field, multiplied by the average time, \(\tau\), that has elapsed since the last collision \(u = \frac{eE}{m}\). This extra velocity appears macroscopically as the drift velocity. In the next encounter, the extra energy, on the average, is lost in the collision through recoil or excitation. Therefore there is a balance between the energy picked up and the collision losses (for more details see Ref. [2]).
This is the result of a classical argument and is known as the Einstein relation.

Unlike positive ions, the mobility for electrons is much greater and is found to be a function of $E$. The electrons take advantage of their small mass and increase their velocity to high values:

$$u = \frac{e E \tau}{2 m_e},$$

where $\tau$ is the mean time between two collisions of an electron ($e, m_e$) in an electric field $E$.

Velocities as high as a few times $10^6$ cm/s can generally be attained before saturation sets in. The electric fields at this point are generally on the order of $1 \text{kV/cm} - \text{atm}$. Figure 3.7 shows some measured results for electrons in different gas mixtures.

The gain in velocity of the electrons may also affect the diffusion rate if the mean energy of the electrons exceeds thermal energies. The factor $k T$ in (3.26) is then replaced by this mean energy, the diffusion constant $D$ then increases accordingly causing a greater
spread of the electron cloud as given by (3.20) and (3.21). This has important consequences for detectors such as the drift chambers which attempt to determine the position of a track by measuring the drift time of the ionization electrons.

Figure 3.7. Drift Velocity as a function of electric field for several gas mixtures by GARFIELD simulation program.
3.4. Avalanche Multiplication

Multiplication\(^4\) in gas detectors occurs when the primary ionization electrons gain sufficient energy from the accelerating electric field to also ionize gas molecules. The resulting secondary electrons then produce tertiary ionization and so on. This results in the formation of an avalanche. Because of the greater mobility of the electrons, the avalanche has the form of a liquid-drop with the electrons grouped near the head and the slower ions trailing behind as shown in Fig. 3.8.

\[\text{Figure 3.8. Avalanche formation. Since the electrons are more mobile than the positive ions, the avalanche takes on the form of a liquid drop with the electrons at the head.}\]

\(^4\) For the calculation of multiplication see next chapter.
The principle of operation of this innovative detector is explained in this chapter. Its electric field, gain, energy resolution and spatial resolution will be defined, too.

4.1. The Description

MICROMEGAS (MICRO MEsh GAseous Structure – Fig. 4.1) [1] is a miniaturized version of a very asymmetric two stage parallel plate detector. A micromesh separates the conversion space, of about 3 mm, from a small amplification gap that can be as small as 100 μm. This configuration allows us to obtain, by applying reasonable voltages in the three electrodes (drift electrode, micromesh and anode electrode), a very high electric field in the amplification region and a quite low electric field in the drift region. Therefore, the ratio between the electric field in the amplification gap and that in the conversion gap can be tuned to large values, as is required for an optimal functioning of the device. Such a high ratio is also required in order to catch the ions in the small amplification gap: under the action of the high electric field, the ion cloud is quickly collected on the micromesh and only a small part of it, inversely proportional to the electric field ratio, escapes to the conversion region. Fig. 4.2 shows a schematic representation of a typical detector. It consists of the following components: drift
electrode (conversion or drift gap), micromesh, anode electrode (amplification gap), strips, and pillars (or spacers).

**Figure 4.1.** The 3D structure of a simple MICROMEGAS

MICROMEGAS, as we have already mentioned, is a two-stage parallel plate avalanche chamber in which several innovative properties rely on a narrow amplification gap, typically $50–100 \mu m$, between two parallel electrodes, the cathode and anode conducting plates [1]. The cathode is made of a thin metallic micromesh, few microns thick, the anode microelements (strips or pads) of a conductor, printed on a insulator board. The technological challenge in such a detector is to keep the small gap constant over the active area. The cathode–anode distance is kept by small insulating pillars, deposited by standard photographic methods on the anode or cathode, covering a small part (1%) of the surface. This technical solution permitted the construction, at low cost, of large chambers$^5$ up to $40 \times 40 \, \text{cm}^2$ and can be extended to larger surfaces with excellent uniformity (10%) and energy resolution over the whole surface [2].

$^5$ For example, a chamber can have dimensions of 5×5 cm², 15×15 cm², 40×40 cm² or even more.
In most applications, a third electrode is placed parallel to the mesh to define a larger gas-filled region where electrons, released by any conversion process (ionizing particle or photon conversion) are drifted towards the mesh\(^6\). On top of the mesh, the electrons are transferred to the amplification gap where they are multiplied in an avalanche process, resulting in a large number of electron-ion pairs. Detectable signals are induced on the anode elements and on the cathode mesh.

### 4.2. The Operation Principles

When a particle crosses the detector, it will firstly pass through the drift electrode. While it is passing the electrode, it is already inside the conversion region, which stretches up to some mm until the grid (or else micromesh – Fig 4.3)\(^7\). It holds a quite weak electric field –of the order of 1 kV/cm– and is the place where the ion-electron pair production takes place. The role of the grid is multiple, and does more than marking the end of the conversion gap and the beginning of the amplification one [3].

---

\(^6\) The mesh is a metallic grid, made of nickel, obtained by electroforming technique. Recent developments of etching techniques, widely used for other micro-pattern detectors [9], had opened the way to a novel technology for the fabrication of the micromesh that is a key element of the Micromegas detector. This new type of mesh is based on simple chemical-etching techniques on a single foil of kapton copper plated on both sides. The manufacture process relies on the high accuracy of the photolithography technique reached in CERN that allows printing on a 5 mm copper a grid with 25 mm openings and a pitch of 50 mm; the kapton is then removed by etching. In a second fabrication, the kapton is partially removed leaving kapton pillars that are used as spacers for the amplification gap of the detector (see Fig. 4.3). The detailed procedure is as follows (Fig. 4.4) [3]:

- A double-sided kapton foil, 50 mm thick, is stretched on a Stesalite frame.
- A solid photoresist of 15 mm thick is applied on the two faces of the copper clad.
- Two lithographic masks are used for both faces of the kapton: one side with the holes pattern and the other with the pillars pattern.
- Copper and kapton are then etched providing the final mesh and pillars.

\(^7\) The MICROMEGAS detector of CAST Experiment has a conversion gap of 2.5 mm, but those that are going to be used in ATLAS Experiment have gap of 100 \(\mu m\).
The voltage applied to it (up to 500 V) is such that the ratio of the electric field in the amplification gap over the field of the conversion gap is very big. The bigger the ratio the higher the electron transmission to the amplification gap reached (in practice a ratio of 20 means full transmission). Once in the amplification gap, the process of avalanche is easily started; the gap is so small (of the order of 50-100 \( \mu m \)) that the electric field achieved is very high (up to 50 kV/cm). At the same time as providing a smooth way for the electrons into the amplification gap, the micromesh prevents the ions produced by the avalanche to enter the conversion gap [4].

**Figure 4.2.** A schematic view of MICROMEGAS:
the micromesh separates the detector volume to the conversion gap (some mm or \( \mu m \)) and the amplification gap (of the order of 50-100 \( \mu m \)) which ends to the strip plane.

While the ions are collected by the micromesh with a high efficiency and speed, the electrons continue in the amplification gap and end their travel on the anode electrode. The anode electrode consists of copper strips with a typical width of 150 \( \mu m \) and a pitch of 200 \( \mu m \), grounded through low-noise charge preamplifiers of high gain to an isolating layer (usually kapton) [4].
The advantages of the technique introduced with MICROMEGAS are listed below (more details in below sections) [4]:

- The fast response: because of the very small path the ions need to travel (amplification gap length ~ 100 μm) and of the very strong field, the ions are very rapidly collected, suppressing any space-charge effects.

- Any mechanical imperfection on the stretching of the micromesh above the strips is compensated, leading to essentially steady gain; an approximation of the change in the amplification factor $M$ with the amplification gap $d$ is given by

$$
\frac{\delta M}{M} = ad \left(1 - \frac{Bpd}{V}\right) \frac{\delta d}{d}
$$

(4.1)

for pressure $p$, applied voltage $V$ and $B$ a constant depending on the gas used (see §4.1.4). Under constant pressure, when $d$ decreases, the multiplication factor increases up to a maximum (for $d = V/B$) and then decreases for higher values of $d$. The combination of the amplification gap and the applied voltage in the MICROMEGAS detectors is such, that the multiplication factor is maximized, so fluctuations due to defects of flatness between the mesh and the anode plane are canceled.
• Counting capability of the order of $10^6$ counts mm$^2$ s$^{-1}$ due to the fast evacuation of the ions and the high granularity of the mesh.
• Because of the constant field along the amplification region, the signal detected in the anode is equally due to the ions and the electrons, contrary to the wire chambers.
• An excellent spatial resolution.

MICROMEGAS has been used in a variety of domains, including high energy physics (COMPASS, n TOF, NA48, TESLA), non accelerator physics (CAST, HELLAZ) and in medical applications (X-ray imaging), reaching performances like the ones in Table 4.1.

Table 4.1. The performance of MICROMEGAS in various applications.

<table>
<thead>
<tr>
<th>Property</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Spatial resolution</td>
<td>12 $\mu$m (rms) [5]</td>
</tr>
<tr>
<td>Time resolution</td>
<td>0.2 ns (rms) [6]</td>
</tr>
<tr>
<td>Energy resolution (at 5.9 keV)</td>
<td>11% (FWHM) [3]</td>
</tr>
<tr>
<td>Rise time of the fast signal</td>
<td>&lt; 1 ns [7]</td>
</tr>
<tr>
<td>Signal-to-noise for M.I.Ps</td>
<td>&gt; 100 [8]</td>
</tr>
</tbody>
</table>

4.3. Electric Field Configuration

Ionization electrons [1], created by the energy deposition of an incident charged particle in the conversion gap, drift and can be transferred through the cathode micromesh; they are amplified in the small gap, between anode and cathode, under the action of the electric field, which is high in this region. The electron cloud is finally collected by the anode microstrips, while the positive ions are drifting in the opposite direction and are collected on the micromesh. The electric field (Fig. 4.5, 4.6 & 4.7) must be uniform in both conversion and amplification spaces. This is easily obtained by using the micromesh as the middle electrode. The electric-field shape is, however, disturbed close to the holes of the micromesh. The knowledge of the shape of the field lines close to the micromesh is a fundamental issue for the operation of our detector, and especially for the efficiency of
the passage of electrons through the micromesh, as well as, for the fast evacuation of the positive-ion build-up.

Figure 4.5. Map of the electric field lines around the micromesh (50 \( \mu \text{m} \) step, 37 \( \mu \text{m} \) diameter of the openings). No voltage on the strips.

Figure 4.6. The electric field in the detector where strips hold a voltage of 60-100 kV/cm: 3D field simulations (FEM)\(^8\) provide field maps by evaluating complex structures of conductors and dielectrics.

\(^8\) Computation of field maps using 3D Finite Element Method and GARFIELD simulation program.
Figure 4.7. (above) Electric field surface and contours for a typical MICROMEGAS, (below) Electric field along the axis of a mesh hole as distance between the mesh and the anode is varied, by MAXWELL 3D simulation program
In a first approximation the strip structure was neglected assuming a plane electrode. Hence, symmetry considerations lead to a model which is limited on a 1/8 of an elementary cell. The electrostatic potential follows Laplace’s law $\Delta V = 0$. Since temperature follows Laplace’s law, a useful analogy with the electrical potential can be made. Boundary conditions are deduced from the symmetry of the model. As no field-lines are leaving the cell, it is an adiabatic case, and the electrical potential is set on the three electrodes. In the following section (§4.1.3.1) we will give details of these calculations and of how the electron transmission, as well as the ion collection on the micromesh, is defined. The results are given as a function of the parameter $\xi$, which is the ratio between the electric fields in the amplification and in the conversion regions.

Fig. 4.8 describes the evolution of the held-line configuration for various values of the ratio $\xi$. For low values of $\xi$, most of the lines leaving the first electrode reach the micromesh. So, in that case the electron transmission is poor. For large values of $\xi$, which correspond to the regular operation of the detector, most of the lines of the conversion field pass through the hole of the micromesh and reach directly the third anode electrode. In this case we reach a full electron transmission. On the other hand most of the lines on the anode plane end in the micromesh, so ions created during the avalanche process will be collected by the micromesh with a high efficiency. Fig. 4.9 shows the electron and ion transmission through the micromesh. For large values of $\xi$, the electron transparency becomes 1, while the ion transmission is extremely low and becomes lower still for larger values of the ratio $\xi$. We would like to point out that the simulation shows that the optical transparency of the micromesh is not a fundamental issue, since electrons can be transmitted by applying higher field ratio [1].
4.3.1. Electron-Ion Transmission

The model used [1], simulating the electric-field configuration near the micromesh, can be represented with a set of three parameters, \( \xi, \eta \) and \( f \) where \( \xi \) is the ratio of the electric field in the avalanche gap to the field in the conversion gap, \( \eta \) is the optical transparency.
of the micromesh and \( f \) is a form factor set to the ratio of the mesh width to the mesh pitch.

The calculation provides directly the electrostatic potential \( V \) in the cell and the electric field is given by \( E = -g d V \). The flows in the three electrodes are \( \phi_1 \), \( \phi_2 \) and \( \phi_3 \) respectively, they can be deduced by the formula \( \phi = \int E \cdot dS \). In order to find the electronic transparency it is useful to express \( \phi_1 \), \( \phi_2 \) and \( \phi_3 \) in terms of partial flows between two electrodes \( \phi_{13} \), \( \phi_{12} \) and \( \phi_{23} \) as it is illustrated on Fig. 4.10. The quantities are related by

\[
\begin{align*}
\phi_1 &= \phi_{13} + \phi_{12}, \\
\phi_2 &= \phi_{23} - \phi_{12}, \\
\phi_3 &= -\phi_{13} - \phi_{23},
\end{align*}
\]

in agreement with \( \phi_1 + \phi_2 + \phi_3 = 0 \).

\[\text{Figure 4.10. Electric field configuration and partial flows}\]
The electronic transparency is given by

\[ \eta_e = \frac{\phi_{13}}{\phi_{13} + \phi_{12}} , \]  

(4.3)

and is a function of \( \xi, \eta \) and \( f \). The ion transparency, supposing a diffusion process, can be expressed as:

\[ \eta_i = \frac{\phi_{13}}{\phi_{13} + \phi_{23}} . \]  

(4.4)

Symmetry considerations between the two gaps separated by the micromesh lead to the conclusion that \( \eta_i \) and \( \eta_e \) are not independent and can be related by [1]:

\[ \eta_i(\xi, \eta, f) = \eta_e \left( \frac{1}{\xi} \eta, f \right) . \]  

(4.5)

Some laboratory tests had been done by Y. Giomataris et al. [1] in order to measure microgrid electron transmission. In a separate setup they had tested the validity of the simulation model. The electroformed micromesh (HV2 – see Fig. 4.11) was placed between two electrodes at a distance of 5 mm to define two gaps, both operating in the ionization chamber mode. A \( ^{254}\text{Cf} \) radioactive source was placed on top operating in the ionization chamber mode. A \( ^{254}\text{Cf} \) radioactive source was placed on top of one of the electrodes (HV1). Each fission created several millions of primary electrons in the gas mixture, a signal detectable by our high-gain charge preamplifiers. The electron cloud was created near the top electrode; through the action of the electric field applied in this gap, it drifted to the middle electrode, the micromesh, giving a signal \( S_1 \). Part of the electron cloud could be transmitted to the second gap when an electric field was applied there, giving a signal \( S_2 \). The transmission was defined as the ratio between the two signals \( \eta = S_1 / S_2 \). They had studied the electron transmission through the micromesh for various electric fields applied on the two gaps. Fig. 4.12 shows the electron transmission as a function of the ratio of the two electric fields: \( \xi = E_2 / E_1 \). There is a fast rise up to a ratio of about 10; then a slower rise up to a ratio 40; finally a plateau when the full electron cloud is transmitted through the micromesh. This is a particular behaviour,
which characterizes the micromesh, and it is different from the measurements obtained with conventional wiremeshes having larger holes.

Figure 4.11. A 3D view of the detector elements

Figure 4.12. Measured electron transmission through the micromesh for values of the $\xi$ ratio

4.3.2. Ion Feedback

As we have already mentioned, MICROMEGAS has the capability to naturally stop most of the ions produced in the amplification space. Now, we are going to define the ion feedback.
Due to the very large field ratio $\zeta$ [12] between the multiplication and the drift regions (as high as 400 or 500) the electric field lines are very much compressed between the two regions (‘‘funnel’’ effect). Following the Gauss theorem, the compression factor of the field lines is equal to the field ratio $\zeta$. But due to collisions in the gas, electrons do not drift along the field lines. They diffuse, especially also in the multiplication space: the transverse extension $s$ (standard deviation) of the avalanche due to diffusion is of the order of $10\sim15 \mu m$; depending on the gas mixture, the electric field and the gap width; in Fig. 4.13 is shown a simulation by GARFIELD of electron diffusion and multiplication in the drift and the multiplication gaps (see also Fig. 4.14). This electron cloud size is much larger than the size of the funnel end (1\sim2 \mu m in radius in the TPC conditions). But ions, due to their high mass, are not submitted to diffusion and drift along the field lines. Assuming that they are emitted with the same distribution as the avalanche, most of them are naturally collected by the micromesh (see Fig. 4.15), and only the fraction of ions created in the small funnel will flow back into the drift volume.

Analytic calculations [12] have been done assuming a bi-dimensional Gaussian distribution of the electron diffusion in the multiplication space. Then it is assumed that ions are emitted from the anode plane, with the same Gaussian distribution (rms $\sigma$) as the avalanche. This is valid since the gain is generally large enough (at least a few hundred), and most of the ions are emitted at a very small distance (a few $\mu m$ only) from the anode plane. Then ions are supposed to drift along the field lines without any diffusion, because they are heavy. Ions emitted outside the funnel will follow field lines ending on the micromesh, and will be naturally collected by it; a very small fraction, produced in the thin funnel, will drift along field lines flowing from the drift volume, and will feed it, before being collected by the HV electrode of the TPC after a very long time (typically a few hundred ms for a 2m drift length).

Following the previous assumptions, it is easy to compute the ion backflow fraction $b$ as a function of the field ratio $\zeta$: As expected the key parameter is the relative value of the size of the ion cloud ($\sigma$) and the mesh pitch ($l$). In Fig. 4.16 the product $\zeta \cdot b$ as a function
of $\sigma/l$ is shown: if this parameter is small (small diffusion and/or too small pitch mesh) ion feedback $b$ is substantially larger than the inverse of the field ratio $\xi$; if $\sigma/l$ is greater than 0.5, the optimal is reached, with an ion feedback equal to $1/\xi$.

Figure 4.13. Left: GARFIELD simulation of electron drift and multiplication in MICROMEGAS. Large ionization corresponds to production of $\delta$-electrons because of passing a muon in a gas mixture Ar + CO$_2$ (70:30) at temperature of 300 $K$ and pressure of 1 atm. Right: schematic view of MICROMEGAS

Figure 4.14. Electron drift lines from a track by GARFIELD simulation program, on the left 20 equal spaced point particles pass through the detector (in Ar + CO$_2$ (80:20)) while on the right a proton of 10 MeV passes the detector (in Ar + iC$_4$H$_{10}$ (90:10)) at temperature of 300 $K$ and pressure of 1 atm
This condition is easily reached: for most of the usual gas mixtures, the transverse diffusion at high electric field (40 – 70 kV/cm) is of the order of \(120 – 150 \mu m \text{ cm}^{-1}\) i.e. \(\sigma = 12 – 15 \mu m\) for a 100 \(\mu m\) amplification gap. With a 500 lpi (lines per inch) micromesh (50 \(\mu m\) pitch), \(\sigma/\ell\) is equal to 0.25 – 0.3, and the ion backflow is 2 or 3 times larger than the optimal value \(1/\xi\). With a 1000 or 1500 lpi mesh (25 or 17 \(\mu m\) pitch) \(\sigma/\ell\) is larger than 0.5, and the expected feedback is equal to the inverse of the field ratio \(\xi\).

As a conclusion it is expected that the optimal ion backflow conditions will be fulfilled with a 1000 lpi mesh for 100 \(\mu m\) gap, and with a 1500 lpi mesh for 50 \(\mu m\). In addition, since ions due to their mass, are not very sensitive to magnetic field, it is expected that ion backflow will not be affected by its presence.

**Figure 4.15.** The field lines in MICROMEGAS by GARFIELD simulation program, Left: Drift lines from an electron track passing through MICROMEGAS with gas mixture Ar + iC\(_4\)H\(_{10}\) (90:10) at pressure of 1 atm and temperature of 300 K. Right: Electron drift lines from a track of 50 equaled spaced points passing through MICROMEGAS with gas mixture Ar + CO\(_2\) (80:20) at pressure of 1 atm and temperature of 300 K.
In order to identify that the theoretical calculations for the ion feedback is met the experimental results Colas et al. [12] make the following experiments: An intensive ($10 mA - 10 keV$) X-ray gun produced primary electrons (see Fig. 4.17) in the 3 mm drift space. The Ni micro-mesh, manufactured at CERN, was located at a distance of 100 $\mu m$ from the anode plane, and the typical gain was a few hundred. Gas mixtures were Argon with 10% Isobutane or 2–3% CH₄. Currents on the drift ($i_d$) and mesh ($i_m$) electrodes were accurately measured. The primary ionisation current $i_p$, which is of the order of a few 10 pA, was obtained by measuring the drift current without gain (by lowering the voltage on the mesh). From these current measurements, it is easy to determine the ion backflow $b$:  

$$b = \frac{i_d - i_p}{i_d + i_p}$$  

as a function of the field ratio by changing the voltage on the drift electrode ($\xi$ was varying in a large dynamic range, between 10 and 700).

In Fig. 4.18 are shown measurements performed with a 500 lpi electroformed Ni mesh: as expected from calculations, the extension of the avalanche is not large enough as compared with the mesh aperture, and ion backflow is degraded by quite a large factor ($\sim 4$) as compared to $1/\xi$: Then measurements (see Fig. 4.19) have been done with a smaller pitch Ni mesh (1500 lpi, 17 $\mu m$): as expected, the backflow is exactly equal to the inverse of the field ratio over a very large range of field ratios. Finally, measurements have also been done in a superconducting coil, varying the magnetic field from 0 to 2 $T$, without any change in the ion feedback, as expected (see Fig. 4.20).

![Figure 4.16. Computed value of feedback × field ratio ($\xi b$) as a function of transverse diffusion / mesh pitch ($\sigma / l$)](image-url)
Consequently, it has been proven and explained that in a MICROMEGAS device the ion backflow is equal to the inverse of the field ratio between the amplification and the drift.
electric fields, with only a few restrictions on the gas mixture, and on the mesh which should have a small pitch \((<25\,\mu m)\).

### 4.4. Advantage of the Small Gap

An interesting property of MICROMEGAS [2] is that, thanks to its narrow gap, locally small variations of the amplification gap, due to, for instance, mechanical defects, do not induce gain fluctuation; they are compensated by an inverse variation of the amplification coefficient. This behavior can be explained by a simple theory:

The electron multiplication \((M)\) in the uniform electric field between two parallel plates in a gas at a pressure \(p\), is described by:

\[
M = e^{a \cdot d}, \tag{4.6}
\]

where \(d\) is the gap of the two parallel electrodes and \(a\) is the 1st Townsend coefficient, which represents the mean free path of the electron between two ionizations. A good approximation of this coefficient is given by Rose and Korff formula (see Fig. 4.21):

\[
a = p A e^{-B \cdot p/E}, \tag{4.7}
\]

where \(E\) is the electric field and \(A, B\) are parameters depending on the gas mixture.

At high electric field values of the 1st Townsend coefficient saturates because its value approaches the mean free path given by the inelastic collision cross-section. The electric field is \(E=V/d\), where \(V\) is the applied voltage. By substituting Eq. (1.7) in (1.6) we get:

\[
\ln(M) = A \, p \, d \exp\left(-\frac{Bpd}{V}\right). \tag{4.8}
\]

The multiplication factor \(M\) is a function of the quantity \(p \cdot d\). Fig. 4.22 shows \(M\) as a function of the gap \((d)\) for a typical mixture of Ar + 5% DME, and for \(V = 300, 350, 400\) and \(450\,V\) at 1 bar.
As we can see that $M$ rises as $d$ increases, it reaches a maximum and then falls at large values of $d$. The maximum is obtained by a differentiation of the Eq. (1.8), resulting in

$$\frac{\delta M}{M} = ad \left(1 - \frac{Bd}{V}\right) \cdot \frac{\delta d}{d}.$$  

The maximum value is for $d_{\text{max}} = \frac{V}{Bp}$ at $p=1\text{bar}$. The amplification gap chosen in this way depends slightly on the gas mixture; for a given applied potential, the multiplication factor is at maximum in the range of gaps between 30-100 microns. This is the range currently used by the MICROMEGAS detectors. The conclusion is that in this range the multiplication factor is maximized and fluctuations due to defects of flatness of the two parallel electrodes are canceled. In few words such narrow gaps are ideal for an optimal operation of the parallel plate gaseous detectors, since all fluctuations due to mechanical defects, atmospheric pressure or temperature variations are suppressed.

**Figure 4.21.** 1st Townsend Coefficient as function of the electric field: the comparison of the Townsend coefficient between the Rose-Korff approximations (curves) and the Monte-Carlo simulation.
It is quite difficult to verify experimentally the previous calculation, as a large variety of very narrow gaps are needed. It needs to be pointed out that this has been verified for two gaps \((d = 100 \text{ and } 50 \mu m)\), but further work is needed to complete the study. It is much easier to verify the variation of \(M\) with pressure, which is expected to be equivalent to the gap variation.

**Figure 4.22.** Calculated gas gain in Ar + 5% DME as a function of the amplification gap for various potentials applied on the microgrid

Fig. 4.23 shows the multiplication factor obtained from our measurements as a function of the pressure for Ar + 7% cyclohexane at \(V = 270V\) for the gap of 50 microns, while Fig. 4.24 shows the gain in He + 6% Isobutane as function of the gas pressure for various potentials applied on the micro-grid. As we can see, the curve clearly shows that there is a maximum of multiplication at \(p = 500 mbar\). Notice that the optimal operation of a
conventional parallel plate avalanche chamber \((d = 4 \text{ mm})\) is at pressures of the order of 10 mbar.

Figure 4.23. Gain versus pressure for a gas Ar + 7% cyclohexane for a 50 \(\mu\text{m}\) amplification gap. The drift and the micromesh voltage was 1000 and 270 V, respectively.

Figure 4.24. Measured gain in He + 6% Isobutane as function of the gas pressure for various potentials applied on the micro-grid

4.5. Gain Properties

The highest gas gain obtained by a gaseous detector is a key issue for the large number of applications. In particular, a detection of minimum ionizing particles requires a large dynamic range because of the Landau fluctuation of the deposited energy and the emission of heavy ionizing particles. The goal of a “good” detector is to achieve a stable operation before the breakdown, which corresponds to a total charge per avalanche approaching \(10^7 - 10^8\) (so called Rather limit).

MICROMEGAS has been tested with a large variety of gas mixtures. Results have been published for Argon mixtures with various hydrocarbons [8, 10]; the maximum safe gain is close to \(10^5\) with 5-10% addition of Isobutane, and three times higher with a small amount of Cyclohexane. Adding CF\(_4\) to the previous mixtures is important, because it improves the time resolution and the total deposited energy [11]. Neon or He mixtures with hydrocarbons allow an increase of the total charge per single avalanche that
approaches the highest Rather values (about $10^8$). As an example, Fig. 4.25 shows the gas gain measured in He + 6% Isobutane mixture using single photoelectrons produced under UV illumination; the maximum gas gain reached was $\sim 1.8 \times 10^7$.

Mixtures of high-Z gases, such as Krypton or Xenon, are relevant for many applications in X-ray digital radiography, crystallography and synchrotron radiation studies. The general conclusion is that the maximum achievable gas gain increases with heavier hydrocarbon quenchers with lower ionization potentials. As an example, Fig. 4.26 shows the gas gain measured as a function of the applied voltage and for various quenchers added to the Xenon carrier gas. The maximum achievable gas gain is increasing as one goes from the Isobutane ($4500$) to Cyclohexane ($10^4$), and finally to Cyclohexene ($3 \times 10^4$). Such high gas gain gives the required margin factor when a detector has to cope with very-high X-ray environments, or at high-pressure operations.

Figure 4.25. Gas gain measured in He + 6% Isobutane as a function of the applied potential
4.5.1. Gain and Grid Geometry

Gain as a function of grid geometry [15] has been measured for three 75 μm gap InGrids, each having a different hole diameter and hole pitch summarized in Table 4.2.

The gain is clearly dependent on the hole diameter (Fig. 4.27). This is understood by looking at the field strength along the hole axis for different hole diameter, keeping the hole pitch fixed (Fig. 4.28). While decreasing the hole diameter, the electric field along the hole axis is higher over a longer distance, therefore the overall gain increases.

Table 4.2. Measured gain for different InGrid geometries

<table>
<thead>
<tr>
<th>Hole pitch [μm]</th>
<th>InGrid 1</th>
<th>InGrid 2</th>
<th>InGrid 3</th>
</tr>
</thead>
<tbody>
<tr>
<td>Hole diameter [μm]</td>
<td>32</td>
<td>36</td>
<td>75</td>
</tr>
<tr>
<td>Gain at 500 V</td>
<td>9800</td>
<td>9900</td>
<td>2900</td>
</tr>
</tbody>
</table>
Figure 4.27. Gain for three grid geometries (see Table 4.2)

Figure 4.28. Electric field along the hole axis of a 50-mm gap InGrid for various hole diameters (400V on the grid).

In Fig. 4.29 & 4.30 we see the gain as a function of mesh spacing for several gas mixtures. The point where the gain doesn't depend on the precise anode-mesh distance is located very low in Ar + CO$_2$ (80:20) (see Fig. 4.29)! At 100 $\mu$m spacing, a gain of $10^4$ calls for a field $E \sim 65$ kV/cm.

4.5.2. Efficiency

With the gas mixture of 90% argon and 10% DME [7], we observe an efficiency plateau of 25 V (Fig. 4.31) at a level of 99%. Because of a non optimized timing of the multiplexed electronic and the length of cables we have lost a factor 7 in the signal. So in fact, the length of the plateau should be 50 V larger.

A typical distribution of the collected charge is shown in the enclosed Fig. 4.31. The ratio of the charge of the cluster (S) to the noise (N) of a strip is 40 at the peak for 395 V. This shape looks like a Landau distribution as expected by the fluctuation in the number of electrons created in the conversion gap.
Figure 4.29. Gain as a function of Mesh distance for gas mixture Ar + CO₂ (80:20) by GARFIELD simulation program.

Several gas mixtures have been tested and detailed results will be published in a coming paper. We have found a very satisfactory new gas filling: argon + cyclohexane with a possible adjunction of CF₄. At a pressure of 1 bar we can see (Fig. 4.32) the dependence of the maximum gain which can be reached as a function of cyclohexane concentration in argon, with a 50 lm gap. The optimum is around 4% and permits a maximum gain of 3.5×10³. Fig. 4.33 shows the efficiency versus the voltage applied on the micromesh for two different conversion gaps obtained with 10 GeV pions. The plateau is quite large, 50 V at 340 V, with a 3 mm conversion gap. The efficiency of the 1 mm conversion gap reaches 96%, with a narrower plateau of 25 V, but it is expected to be larger for the next scheduled beam-tests, with improved electronics.
The admixture of various proportions of CF$_4$ has a strong influence on both electron drift velocity and electron diffusion in the conversion gap: the electron velocity increases and the diffusion coefficient decreases. With 20% of CF$_4$ admixture the maximum gain is still high, $1.5 \times 10^5$. The most promising mixture is pure CF$_4$ and a small proportion of
cyclohexane. It increases by more than a factor 2 the number of primary electrons created by the incident particle.

Figure 4.31. Efficiency plateau obtained with 10 Gev pions. Enclosed the charge distribution of the collected signal (Landau distribution)

Figure 4.32. Gain as a function of the amplification electric field for different cyclohexane concentrations in argon
As shown in Fig. 4.34, the maximum gain achieved with 7% cyclohexane is high, $6 \times 10^4$, when the detector is irradiated with a 5.9 keV radioactive source, at low flux. It is still significant, $10^4$, at high flux, $5 \times 10^5$ Hz/mm$^2$, in a 8 keV X-ray environment. Filling the detector with that mixture we expect to achieve full efficiency with a conversion gap below 1 mm. Such a small gap improves the time resolution and avoids the deterioration of the space resolution due to inclined tracks or the effect of the magnetic field.

Figure 4.33. Efficiency plateau for 3 and 1mm conversion gap
4.6. Energy Resolution

The resolution of gaseous detectors [15] is mainly determined by the primary charge fluctuations and the single electron gain fluctuations. These sources are intrinsic to the sensing gas and cannot be avoided. On the other hand, fluctuation sources like attachment and collection efficiency depend on the drift field and can therefore be optimized for minimum energy resolution.

Energy resolution is calculated using $^{55}$Fe spectra. $^{55}$Fe emits quanta of 5.9 and 6.5 keV in the ratio 9-1. This ratio is slightly modified to 7.5-1, due to the different absorption of these lines in the gas [16]. Spectra were fitted using two Gaussian functions (Fig. 4.35). The parameters (mean, height, width) of the 6.5 keV line were fixed by the ones of the 5.9 keV line. The energy resolution is defined as the FWHM of the 5.9 keV line.

For the three grid geometries tabulated in Table 4.2, the energy resolution was measured as a function of the grid voltage. Remarkably, the three curves almost superimpose when plotted as a function of the gain (Fig. 4.36).
As already noted in Ref. [3], the resolution exhibits a minimum with respect to grid voltage (or gain). An explanation of the resolution improvement could be the reduction of avalanche fluctuations when increasing the amplification field i.e. transition from exponential to Polya single electron gain fluctuations. Degradation above a gain of $5 \times 10^3$ could be explained by secondary avalanches induced by UV photons or space charge effects that distort the field.

Figure 4.35. $^{55}\text{Fe}$ spectra in Argon showing the escape peak of Argon and the peak of $^{55}\text{Fe}$ (5.9 keV)

Figure 4.36. Resolution vs. gain for various hole diameters
4.7. Studies on Space Charge

At fixed voltages, the efficiency will generally decrease as the incident flux is increased. This is the result of a space charge build-up in the amplification region of the detector due to positive ions from avalanches drifting back through the mesh. Since their mobility is low, an accumulation of ions occurs as the number of avalanches increase. The effect of this charge is to alter the electric field of conversion region. Like the problem of TPC (Time Projection Chamber) which arises during operation is the accumulation of a space charge in the drift volume. The ions, which are produced by the avalanche, are sufficiently numerous that a distortion of the electron field in the drifting volume occurs. Hence, it is necessary to find how the space charge affects the MICROMEGAS detector.

To calculate the space charge in MICROMEGAS we are going to exam the following issues: a) Steady-State theory, and b) Equilibrium State.

4.7.1. Steady-State Theory

In the presence of an electric field, the electrons and ions freed by radiation are accelerated along the field lines towards the anode and cathode respectively. This acceleration is interrupted by collisions with the gas molecules which limit the maximum average velocity which can be attained by the charge along the field direction. The average velocity attained is known as the drift velocity, \( u \), of the charge and is superimposed upon its normal random movement. Compared to their thermal velocities, the drift speed of the ions is slow; however, for electrons this can be much higher since they are much lighter.

In kinetic theory, it is useful to define the mobility (Fig. 4.37 & 4.38) of a charge as

\[
\mu = \frac{u}{E} \quad (4.9)
\]

where \( u \) is the drift velocity and \( E \) the electric field strength. For positive ions, the drift velocity is found to depend linearly on the ratio \( E/p \), (also known as the reduced electric field), up to relatively high electric fields. At a constant pressure, this implies that
the mobility $\mu$ is a constant. For a given $E$, it is also quite clear that $\mu$ varies as the inverse of the pressure $p$.

**Figure 4.37.** Mobility of Argon and CO$_2$ ions in Ar as a function of Electric field for a temperature of 300 K by GARFIELD simulation program

**Figure 4.38.** Mobility of Neon and CO$_2$ ions in Ne as a function of Electric field for a temperature of 300 K by GARFIELD simulation program

The steady state theory requires that new matter must be continuously created (mostly as hydrogen) to keep the average density of matter equal over time. Hence, if that assumption is true, then the quantity of $\rho \cdot u$ is constant, where $\rho$ is the density, and then from the Eq. (4.9)

$$\rho(y) = \frac{\sigma}{\mu \cdot E(y)},$$

(4.10)

where $\sigma$ is the number of ions $cm^2\cdot\mu s$ produced by a particle passing through the conversion region. From the 1$^{st}$ Maxwell equation, Gauss’ Law,

$$\rho(y) = \frac{\sigma}{\mu \cdot E(y)},$$

(4.11)

where $\int E \cdot d \bar{s} = \int [E(y) - E(mesh)] \cdot dy$, $\frac{1}{\varepsilon_0} \int \rho \cdot dV = \frac{1}{\varepsilon_0} \int_0^y \frac{\sigma}{\mu \cdot E(y)} \cdot dy$, and $\varepsilon_0$ is the vacuum permittivity, we have
\[ E(y) = E(mesh) + \frac{1}{\varepsilon_0} \int_{mesh}^{y} \frac{\sigma}{\mu} \cdot E(y') \cdot dy'. \] (4.12)

and after the derivation of the above equation and then its integration we finally take the following expression

\[ E(y) = \sqrt{E_0^2 + \frac{2 \cdot \sigma}{\mu} \cdot y}, \] (4.13)

where \( E_0^2 \) is an integration constant which follows the bounder condition

\[ \Delta V = \int_{0}^{y_{gap}} E(y) \cdot dy. \]

Consequently, Fig. 4.39 shows the electric field as a function of the length \( y \) due to different rates of muons and for \( \sigma / \varepsilon_0 = 3.6 \cdot 10^{-5} \text{ ions} \cdot V / (\text{cm} \cdot \mu s) \), \( \mu \approx 1.5 \cdot 10^{-6} \text{ cm}^2 / (\text{V} \cdot \mu s) \), \( \Delta V = 600 V \) and \( \text{gap} = 3 \text{ mm} \), while Fig. 4.40 shows the potential as a function of the length \( y \) due to different rates of muons.

**Figure 4.39.** Electric field as a function of the length \( y \) for different rates of muons by GARFIELD simulation program

**Figure 4.40.** Potential as a function of the length \( y \) for different rates of muons by GARFIELD simulation program
4.7.2. Equilibrium State

In order to avoid field distortions [13] in the conversion region, the amount of back drifting ions from the gas gain region towards the drift cathode should be kept as small as possible, again at least for high-rate applications. To estimate the maximum acceptable ion feedback, the following ion tube model is applied describing space charge effects in the drift region: a continuous point like photon beam produces after a short period an equilibrium state where the positive charged ions move within an ion tube from the anode to the drift cathode (see Fig. 4.41). Since the distance between the anode and the gas gain structure is negligibly small compared to the drift length of usually several millimeters only the conversion region between the upper side of the mesh and the drift cathode is further considered. Averaged over a larger period of time the total charge of

\[
Q_{\text{total}} = (\overline{t}_a + t_{d_{\text{max}}}) G b \varepsilon N_{\text{prim}} R q_e
\]  

(4.14)

is drifting towards the drift cathode. \(N_{\text{prim}}\) denotes the # of ions per photon produced in the conversion region, \(\overline{t}_a\) is the mean drift time of these ions (depending on the absorption coefficient of the gas, the \(\mu\) energy and the maximum drift length), \(t_{d_{\text{max}}}\) is the ion drift time for the way from the mesh to the drift cathode, \(G\) is the effective gain, \(b\) is the ion feedback, \(\varepsilon\) is the absolute electron transparency, \(R\) is the incoming muon rate, and \(q_e\) is the electron charge. The second term of the Eq. (4.14) becomes dominant for a large electron transparency, an ion feedback in the range of some percent and a gas gain larger than about 1000:

\[
Q_{\text{con}} \approx N_{\text{prim}} R q_e t_{d_{\text{max}}} G b \varepsilon
\]  

(4.15)

The ion drift time \(t_{d_{\text{max}}}\) can be calculated as the quotient of the total drift length \(l\) and the ion drift velocity \(u_{\text{ion}} = \frac{\mu_{\text{ion}} \cdot E_1}{p}\)

\[
t_{d_{\text{max}}} = \frac{l p}{\mu_{\text{ion}} E_1},
\]  

(4.16)
where $\mu_{\text{ion}}$ is the ion mobility, $E_1$ electric field of conversion region, and $p$ the gas pressure.

Using Gauss’ law, the electric field $E$ produced by the ion charge cloud can be calculated

$$
\oint_s \vec{E} \cdot d\vec{s} = \frac{Q_{\text{ion}}}{\varepsilon_0},
$$

(4.17)

where $s$ is the surface of ion cylinder, and $\varepsilon_0$ the vacuum permittivity. The length of the tube $l$ is given by the distance between the upper side of the mesh and the drift cathode; the radius $r_0$ is in a first approximation determined by the transverse diffusion of the primary electrons in the conversion region. Already for drift lengths of some millimeters one obtains, that $r_0 \ll l$. Therefore, the area of the ion tube’s front sides ($\propto r_0^2$) are small compared to the surface of the cylinder ($\propto r_0 l$), hence the electric field has only a radial component $E_r$. Solving the integral in Eq. (4.17) and dividing it by $E_1$ the ratio between the radial electric field component $E_r$ and the drift field $E_1$ is determined by

$$
\frac{E_r}{E_1} = \frac{Q_{\text{ion}}}{2\pi \varepsilon_0 l} \frac{E_1}{E_1} = \frac{N_{\text{prim}} q_e b \varepsilon G R p}{2\pi \mu_{\text{ion}} E_1^2 r}, \text{ for } r \gg r_0.
$$

(4.18)

**Figure 4.41.** Schematic of the ion tube model.

The back drifting ions with charge $Q_{\text{ion}}$, coming from the gas gain region through the holes of the mesh, move with the velocity $v_{\text{ion}}$ towards the drift cathode within a tube with radius $r_0$. To calculate the effect of this space charge, a weak spot is placed close to the hot spot with a distance of $r_S$. 
The expected spatial distortion $r_d$ due to this field perturbation can then be calculated as

$$r_d (r) = \frac{E_r (r)}{E_i} \cdot l$$  \hspace{1cm} (4.19)$$

and increases with low ion mobilities, high pressure and gas gain, rate, muon energy and drift length, whereas it is getting small for huge drift fields and large distances $r$ to the centre of the ion tube.

The electrons, produced in another photon spot, which is placed close to this hot spot at a distance $r_S$, are attracted by the ion tube generated by the strongly illuminated point (see Fig. 4.41). Supposing a moderate gas gain of $G = 2000$, a high spot rate of $R = 10^6 Hz$ of 8 keV photons in 1 bar Ar/CO$_2$ (90:10), respectively $N_{prim} = 300$ and $\mu_{ion} = 1.67 \cdot 10^{-4} m^2 bar/(V \cdot s)$, a drift field of $E_i = 1kV/cm$ and an electron transparency of $\varepsilon = 0.8$ leads to a field ratio $E_r (r)/E_{drift} = 8.3 \cdot 10^{-4} mb/r$. Assuming a two spot configuration with a distance of $r = r_S = 2 mm$, a total drift length of $l = 10 mm$ and a maximum spatial distortion limit of $r_d = 200 \mu m$ due to space charge effects the ion feedback should not exceed $b_{max} = 5\%$ as can be calculated from Eq. (4.24). For higher gas pressures or heavy noble gases with slow ion mobilities like xenon (see Fig. 4.39 & 4.40) the maximum ion feedback should be far below 5\%.

4.8. Signal development and electronics

The signal [2, 7] induced on the anode strips is a sum of the electron and ion signal. The charge induced by electrons and ions are equal. But, with respect to the ions, the drift velocity of the electrons is 100 times bigger, so the current pulse is 100 times shorter and 100 times higher. Using a low-noise charge preamplifier, the charge signal is mainly due to the positive ion drift to the micromesh electrode, which takes place typically within 100 ns, depending on the width of the amplification gap. Figure 4.42 shows signals coming out from a charge-sensitive preamplifier for an Argon + Isobutane (90:10) gas mixture. A reduction of the amplification gap from 100 to 50 \mu m reduces the signal rise
time by a factor 3. Using a gap of 30 $\mu m$ the rise time is only 17 $ns$, another reduction of a factor of two. So, in the latter case, shaping of the signal around 17 $ns$ allows to catch-up the fully induced charge and therefore permits a comfortable operation of the detector at moderate gains. One must also take into consideration that the ion collection time decreases by using higher ion mobility carrier gas (Ne or He).

![Figure 4.42](image)

**Figure 4.42.** Signal given by a charged preamplifier for various amplification gaps: 30, 50, and 100 $\mu m$.

The conclusion is that MICROMEGAS can be used with low-noise charge preamplifiers without loss due to ballistic deficit, which occur in other micro-strip devices; choosing the right amplification and the right gas mixture, the rise of the detector can be compatible with the shaping of the charge amplifier. Due to the faster drift velocity, the electron current is larger and faster (about 1 $ns$ instead of 100 $ns$ for the ion signal). Therefore a very fast rise of the signal, followed by a tail due to the ion drift, is expected. Such fast electron signal is quite difficult to catch, but is within reach with present electronics. For example, using the current-sensitive preamplifier with a fast rise time ($t < 1ns$), the result is spectacular (see Fig. 4.43).

The fast signal has a rise of 1ns and amplitude ten times higher than the ion tail. Such fast signals will allow the development of novel drift chambers or small TCP’s with a time resolution below 1 $ns$. 

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Figure 4.43. Fast electron signal from the current preamplifier in Ar + Isobutane (90:10) gas mixture. Notice that the electron signal and ion tail are developed within 5 ns and 80 ns, respectively.

The time resolution of MICROMEGAS was first investigated using the Lecroy-MQS104 preamplifier [2]. The measurement showed that the time accuracy is dominated by the time-jitter of the ionization clusters produced by the minimum ionizing particles traversing the conversion gap. The best result (4.5 ns) was obtained with a high drift velocity gas: a mixture of Argon, CF$_4$ and Isobutane. It is therefore quite logical to speculate that using CF$_4$ as carrier gas, a conversion depth of 1 mm and faster electronics can reach a time accuracy of 1 ns.

4.8.1. Electron and Ion Signal

The vast majority of the electrons and ions are produced in the immediate vicinity of the pad plane. The avalanche electrons travel only a small distance before being absorbed while the avalanche ions will mostly be absorbed by the mesh. Therefore:

- initial part is happened due to the electrons,
- most of the electron signal are killed by electronics,
• tail of the signal is due to ions, and
• bulk of the integrated avalanche signal is due to ions.

The plots in Fig. 4.44 show the weighting field which has to be multiplied with the electron velocity in order to get the currents for various locations in $z$, where $z$ is perpendicular to the mesh in these plots. The left plot is along a line through the hole, the right plot is along a line through mesh wires. They show that the current in the anode is close to zero for electrons outside the amplification region ($z > 0cm$).

The plots in Fig. 4.45 show the signal shapes without electronics. The ion tail is boring since the weighting field for the complete anode plane is constant and so is the velocity. Hence, this current is constant too. For a single strip, it is a bit different: the weighting field would no longer be constant. The electron current is not constant because the number of electrons grows during the avalanche.

**Figure 4.44.** The weighting field for various locations in $z$, by GARFIELD simulation program. The left plot is along a line through the hole, the right plot is along a line through mesh wires
**Figure 4.45.** The signal shapes of electrons and ions in a gap of 100 μm and at 600 V (no electronics), by GARFIELD simulation program

### 4.8.2. Response of a Strip

In Fig. 4.46 we see the effect of strip width without electronics, while in Fig. 4.47 the signals have now been convoluted with some response function, which is not at all what we are using. This plot is therefore not relevant.

The response of neighbouring pads depends strongly on the electronics transfer function and the time at which the signal is sampled. For a 100 μm gap and 400 μm wide strips:

- A single electron will not give a 3-strip cluster with sizeable signals in all 3 strips.
- Signals will as a rule be seen in 2 adjacent strips.
- Charged particles can give 3-strip clusters through diffusion.
Figure 4.46. The effect of strip width (of 400 $\mu$m) without electronics, by GARFIELD simulation program.

Figure 4.47. The effect of strip width (of 400 $\mu$m) with some electronics, by GARFIELD simulation program.
4.9. Spatial Resolution

In an experiment in 2001, Derré et al. [5] had concentrated on the ultimate space resolution which can be obtained with MICROMEGAS detectors at atmospheric pressure. Hence, in that section the experimental setup is described. Moreover it presents the data analysis and the experimental results which are discussed and compared with simulation.

4.9.1. Experimental Setup Description

MICROMEGAS detector consists of a conversion and drift space which is followed by a narrow amplification gap defined by a cathode plane and an anode plane. The cathode is a Ni-electroformed micromesh with square holes of $39 \times 39 \, \mu m^2$ at a 50.8 $\mu m$ pitch (500 lpi). The anode is a simple PCB with copper strips on an epoxy substrate. The distance between the strips and the micromesh is precisely maintained by small spacers printed on the anode strips using conventional lithography of a photoresistive polyamide film. All detectors they used for this study have a conversion gap of 3 mm, an amplification gap of 100 $\mu m$ and a 1.6 mm thick epoxy substrate.

To measure the spatial resolution a set of MICROMEGAS detectors were installed along a 10 GeV/c pions beam on the T9 line of the CERN PS accelerator. The particles were crossing the detection planes perpendicularly and the divergence of the beam was of the order of 1 mrad. All detector planes have strips parallel to each other, so only one coordinate ($y$) of the track was measured. No information was available in the $x$ direction. To reduce systematic errors due to possible misalignment of the strips (i.e. any defect in the $y$ parallelism), the range along the unknown coordinate $x$ was limited by a trigger scintillator of 2 cm wide in the $x$ direction. For each detector plane the amount of matter, in the active area, corresponds to about 1% of radiation length and was dominated by the PCB epoxy thickness. At 10 GeV/c momentum, the mean deviation angle of the track due to multiple scattering through one detector was about 0.2 mrad. With a 1 mrad beam divergence and this amount of multiple scattering, an attempt to measure the spatial resolution of the studied chamber placed in between the others would fail. Large errors
on the spatial resolution measurement could not be avoided: with a long lever arm the multiple scattering would dominate and with a short lever arm the error on the track angle would become too big. It is worth mentioning here that in experiments where a high-accuracy tracking was required, the amount of material could be reduced if a thin epoxy sheet (down to 100 $\mu m$) or a kapton foil was glued on top of a honeycomb light structure.

To overcome this difficulty they had adopted a doublet configuration. Two MICROMEGAS detectors were mounted back to back at a distance of 10 $mm$. To provide an unambiguous reconstruction of the track and a precise measurement of its angle, at least two other detectors had been used with a total arm length of about half a meter. The spatial resolution was determined from the difference of the positions (centroid of the cluster of adjacent hit strips) measured in each plane of the doublet which is corrected by the track angle. The error on this difference due to multiple scattering was of the order of 2 $\mu m$. As they demonstrated later on, the measurement error in MICROMEGAS never exceeds 50 $\mu m$. So, the error due to the track angle correction is less than 1 $\mu m$. These two sources of error will be neglected in the determination of the accuracy of the spatial resolution.

The chambers they used to track the incident particles have an active area of $14 \times 12 \, cm^2$ with a strip pitch of 317.5 $\mu m$ (70 $\mu m$ inter-strip space), leading to a reasonable number of 384 read-out strips. The spacers are small cylinders of 200 $\mu m$ diameter, 100 $\mu m$ high and set every 2 $mm$.

The doublet of chambers used to investigate the spatial resolution has a strip pitch of 100 $\mu m$ (50 $\mu m$ inter-strip space). However, the active surface was reduced to $5 \times 4 \, cm^2$ in order to keep the same number of read-out channels. The spacers have an oval shape, 100 $\mu m$ long and 50 $\mu m$ large, to fit the strip width. They are disposed every 600 $\mu m$ along every sixth strips. The spacers cover about 17% of the corresponding strip or about 3% of the total active area. This is a safe configuration, although not optimized with respect to dead space.
One early measurement was also performed with a special doublet with strips deposited on glass at a pitch of 50 \( \mu m \) with smaller active area and the same number of read-out channels.

### 4.9.2. Analysis and Results

The first step of the analysis consists in determining for each channel the pedestal distribution. This is done after subtraction for each Gassiplex card of the correlated noise. For the 100 \( \mu m \) doublet used for the spatial resolution measurement, the average rms width of the pedestal distributions is equal to 5 ADC units, equivalent to 1000 electrons (ENC) as expected. Then for each plane, clusters of adjacent channels with a signal above 3 pedestal standard deviations are formed. Dead channel are identified and listed.

To study the efficiency of one detector of the doublet, the incident track is reconstructed in the other chambers. When its intercept with the studied detector is close to a dead region, the track event is abandoned, otherwise the track is used to determine the efficiency. The efficiency is defined as the fraction of tracks for which the residual between the extrapolated intercept and the closest cluster position is less than 0.5 mm. The efficiency depends on the operating conditions, i.e. the voltage set on the micromesh. But even at high gain it also depends on the position of the track with respect to the spacers. Fig. 4.48 (a) and (b) shows typical efficiency responses of each doublet detector as a function of the position (y) of the track intercept. The local loss, at a level of 20% is seen at every sixth strip, corresponds to the strips which hold the spacers and is compatible with the area of the spacers. Later on such an area around a spacer will be considered as a dead space. Note that this affects less than 3% of the total active area. This effect could have been reduced by at least a factor 2 by reducing the size of the spacers.

The detection efficiency increases with the voltage set on the micromesh until a value close to 100% is reached. This happens when the most probable value of the signal-to-
noise ratio is larger than 15. The signal-to-noise ratio \( S / N \) is here defined as the total collected cluster charge \( S \) divided by the mean strip noise \( N \).

![Figure 4.48](image)

**Figure 4.48.** The detection efficiency of (a) Det 1 and (b) Det 2, the two detectors of a MICROMEGAS doublet, as a function of the track intercept \( y \).

### 4.9.2.1. The CF\(_4\)–Isobutane gas mixture

Among all gases investigated, the best spatial resolution was obtained with a CF\(_4\) 80\% and iC\(_4\)H\(_{10}\) 20\% as quencher, gas mixture. In such a mixture a 10 GeV/c pion produces an average of 50 charges in the 3mm conversion gap. With a 1000 electrons noise a gain as low as 300 is in principle sufficient to reach full efficiency. Typical distributions of the total cluster charge collected on strips are shown in Fig. 4.49 (a) and (b) for each of the two detectors of the doublet. They are well fitted by a Landau-type distribution. The corresponding cluster size distributions are shown in Fig. 4.49 (c) and (d). They peak around a cluster size of 2 which is not the most suitable situation to get, with the centroid of the cluster, a good estimation of the position of the intercept of the track in the detectors as will be discussed in the next section. Fig. 4.50 (a) shows the distribution of the difference between the two cluster centroids for a track detected in both chambers of the doublet. This difference has been corrected for the track angle as measured by the other chambers. The distribution is well fitted by a gaussian with a standard deviation of 17 \( \mu m \). If one assumes that both chambers of the doublet have the same resolution and that the centroids measurements are not correlated, one can estimate the spatial resolution of one chamber to be: 17 \( \mu m \) divided by \( \sqrt{2} \), which is 12 \( \mu m \). A discussion on this point at the light of our simulation is given in the next section.
Figure 4.49. The charge response, in ADC units, of (a) Det 1 and (b) Det 2. The line is the Landau distribution fit; the points are the result of the simulation. The corresponding cluster size of (c) Det 1 and (d) Det 2.

Although the resolution is quite insensitive to the gain of the chamber, this is not the case for the drift electric field: the resolution improves by 30% when the electric field in the drift gap increases from 0.5 to 1.5 kV/cm. To get the best resolution, a field of at least 1.5 kV/cm has to be applied in the drift space.

This observation is related to the behavior of the transverse diffusion in the gas. A similar effect has been observed [11] for the longitudinal diffusion.
4.9.2.2. Other gas mixtures

The spatial resolution of MICROMEGAS has also been studied for other gas mixtures. For comparison with the above results, the corrected position difference distribution in a doublet operated with our standard gas mixture Ar + isobutane (90:10) is shown in Fig. 4.50 (b). The measurement shows a clear degradation of the accuracy which is 42.5 μm (standard deviation divided by $\sqrt{2}$).

![Figure 4.50](image)

**Figure 4.50.** The difference between the intercept of tracks in the two detectors of a MICROMEGAS doublet for a gas mixture of (a) CF₄ and isobutane, (b) argon and isobutane. The fitted lines are gaussian distributions.

Table 4.3 summarizes all the results obtained with a doublet of detectors. For each gas mixture we give the measured resolution and the transverse diffusion coefficient. The transverse diffusion coefficient results from an adjustment in our simulation to reproduce the charge and cluster size distributions shown by the data.

The spatial resolution improves when the amount of quencher is increased in the gas mixture and the amount of CF₄ is large. The best result was obtained with CF₄ as carrier.

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gas. With isobutene as quencher the mixture combines a low transverse coefficient (as long as the drift field is higher than 1.5 kV/cm) and a high primary ionization.

Table 4.3. MICROMEGAS spatial resolution

<table>
<thead>
<tr>
<th>Gas Mixtures</th>
<th>Measured Resolution (μm)</th>
<th>Transverse Diffusion Coefficient (μm/√cm)</th>
<th>Strip Pitch (μm)</th>
<th>Drift Field (kV/cm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ar + iC₄H₁₀ (90:10)</td>
<td>42.5</td>
<td>370</td>
<td>100</td>
<td>1.0</td>
</tr>
<tr>
<td>He + iC₄H₁₀ + CF₄ (89:6:5)</td>
<td>35</td>
<td>180</td>
<td>100</td>
<td>1.8</td>
</tr>
<tr>
<td>He + iC₄H₁₀ + CF₄ (84:6:10)</td>
<td>30</td>
<td>160</td>
<td>100</td>
<td>1.8</td>
</tr>
<tr>
<td>He + DME (80:20)</td>
<td>25</td>
<td>130</td>
<td>50</td>
<td>1.0</td>
</tr>
<tr>
<td>CF₄ + iC₄H₁₀ (80:20)</td>
<td>18</td>
<td>170</td>
<td>100</td>
<td>0.4</td>
</tr>
<tr>
<td>CF₄ + iC₄H₁₀ (80:20)</td>
<td>11</td>
<td>100</td>
<td>100</td>
<td>2.7</td>
</tr>
</tbody>
</table>

4.9.3. Monte-Carlo simulation and discussion

4.9.3.1. The simulation

A detailed Monte-Carlo simulation of the detector has been developed. It includes a description of the electromagnetic interaction of the incident particle with the gas and the production of primary pairs of charges. However, delta rays production and recoiling following nuclear interactions have been neglected. The migration of electrons to the amplification gap, the development of the avalanche and the formation of the induced signal on the strips is fully simulated assuming that all field lines end on the top of strips. Electronic noise is added to the charges on the strips. The Monte-Carlo simulation contains two free parameters: the amplification gain and the transverse diffusion coefficient. These are adjusted to reproduce, respectively, the charge and cluster size data distributions. The agreement with the data which is shown in Fig. 4.49, for our best measurement with CF₄ – Isobutane is quite good.
Three effects contribute to the spatial resolution: (1) the transverse diffusion in the gas for which the adjusted coefficient is \(100 \frac{\mu m}{\sqrt{cm}}\); its effect on the resolution is equal to \(8.0 \mu m\), which is given by the mean dispersion of the electrons in the 3 mm drift space (56.4 \(\mu m\)) divided by the square root of their mean number (50); (2) the pitch of the strips inducing a error in the calculation of the centroid of cluster which is estimated by our Monte-Carlo program to be equal to 7.5 \(\mu m\). The quadratic sum of these two effects is equal to 11.0 \(\mu m\); (3) the pitch of the mesh, an effect which was not foreseen before the measurements. It was demonstrated by the simulation and will be discussed below. Its consequence is that the resolution in one chamber depends on the position of the track as long as the relative position of the mesh with respect to the strips is not known.

4.9.3.2. Systematic effects due to the pitch of the mesh

The following discussion is an analytical interpretation of the simulation. First note that the pitch of the mesh holes, 50.8 \(\mu m\), is close to one half of the 100 \(\mu m\) strips pitch. Let us call \(p_s\), the pitch of the strips and \(p_g\) the pitch of the mesh holes. The simulation showed that the current which is induced on the strips by the moving charges behaves as if there was no space between each strips. We, therefore, assume in the following calculation that the strip width is equal to the strip pitch \(p_s\), as if each strip could touch its neighbors.

We then define \(\delta = y_{mesh} - y_{strip}\), the distance between the center of a hole and the beginning of the nearest strip (see Fig. 4.51). Because of the periodicity of the basic cell, consisting of a strip and two mesh holes, we restrict our computation to \(0 \leq \delta \leq p_s\). We should now stress the fact that because of the electric field configuration around the holes, drift electrons produced in front of a given hole will then be collected at the center of this hole. If the strips and the holes are aligned (i.e. \(\delta = p_g / 2\)), this does not matter: a given strip will collect all the drift charges that are really produced in front of it. But if not, electrons that are produced between the end of the strip we consider and the end of the hole (between \(\delta\) and \(p_g / 2\)), will also be collected on this strip, although they were produced in a region lying in front of the next strip. As a consequence, the measured position of the track that produced such electrons will be shifted by a quantity equal to \(p_s\).
If tracks within one cell are uniformly distributed along the y-coordinate, the fraction of such shifted track measurements among all tracks is:

\[ f = \frac{\varepsilon \left( \frac{p_y}{2} - \delta \right)}{p_y}, \]  
(4.20)

where \( \varepsilon \) is the sign of \( \frac{p_y}{2} - \delta \), and the mean shift of the track coordinate is given by

\[ y_{\text{meas}} - y_{\text{track}} = (y_{\text{track}} + \varepsilon p_y) f + y_{\text{track}} (1 - f) - y_{\text{track}} = \frac{p_y}{2} - \delta. \]  
(4.21)

![Figure 4.51. The basic detector layout for the simulation.](image)

Now, in addition, we have to take into account the smearing due to the extension \( \sigma \) of the avalanche. This extension is due to the transverse diffusion in the gas. Our simulation shows that the barycenter of the avalanche is close to \( \frac{1}{4} \) of the amplification gap above the strip plane and the extension is equivalent to a diffusion on a path of 25 \( \mu m \). That gives a standard deviation \( \sigma \) equal to 5 \( \mu m \).

As \( \sigma \) is small with respect to \( p_y \), we can neglect the smearing at a distance from a hole greater than \( p_y \), and the previous equation becomes

\[ y_{\text{meas}} - y_{\text{track}} = \frac{p_y}{2} - \delta - p_y \frac{\text{CL}(\delta)}{2} + p_y \frac{\text{CL}(\delta)}{2} \]  
(4.22)
where \( \frac{CL(x)}{2} = \frac{1}{\sqrt{2\pi} \sigma} \int_x^{\infty} e^{-u^2/2\sigma^2} du \). The two last terms in the previous equation correspond to the smearing due to the avalanche extension issued from the two holes which are in front of a strip. One term is for the hole close to the beginning of the strip, the second for the hole close to the end of the strip. They correspond, respectively, to the mean shift in the track position on the left side (first term) or the right side (second term) for the fraction \( CL(\delta)/2 \) or \( CL(p_2 - \delta)/2 \) of the avalanche which hits on the left or the right neighbor strip.

The result is shown in Fig. 4.52. Note that the mean of the systematic error is null on the total range of \( \delta \). With the 500 lpi mesh we have a systematic error on the position of the track which can be as big as 14 \( \mu m \). The effect is greatly reduced if the pitch of the mesh is smaller by a factor of two (1000 lpi). Since the pitch of the mesh is not exactly a submultiple of the pitch of the strips, \( \delta \) changes from cell to cell. The period of recurrence is 6.35 \( mm \). As we reconstruct tracks in a range of about 2 \( cm \), we can consider \( \delta \) as randomly distributed. In that condition, the systematic mean quadratic error is equal to 9.6 \( \mu m \).

To evaluate the resolution in one detector, we measure the difference of the reconstruction points in a doublet of chambers. In that case we can write

\[
y_{\text{mesh}2} - y_{\text{strip}2} = y_{\text{mesh}1} - y_{\text{strip}1} + (y_{\text{mesh}2} - y_{\text{mesh}1}) - (y_{\text{strip}2} - y_{\text{strip}1}),
\]

or, with the same notation as previously \( \delta' = \delta_1 + \delta_c \) with \( \delta_c = \Delta y_{\text{mesh}} - \Delta y_{\text{strip}} \).

As the two detectors are the same, \( \delta_c \) is a constant, independent of the track position. A shift \( \delta_1 \) in the first detector for a track at \( y_1 \) is associated with a shift \( \delta_2 \) in the second detector at the translated position \( y_1 + \delta_c \). So the shift in the first detector, \( y_{\text{meas}1} - y_{\text{track}} \), for a track at \( y_1 \) corresponds to a shift in the second detector at a position which is translated by a constant. And the difference, \( y_{\text{meas}2} - y_{\text{meas}1} \), is affected by a systematic error which is the difference between these two shifted track positions as shown in Fig.
4.53 (a) for different values of $\delta_c$. By randomizing $\delta_1$ (and the correlated $\delta_2$), we get the expected resolution on the difference as shown in Fig. 4.54.

**Figure 4.52.** The simulated average systematic error on the track intercept in a detector as a function of the position of the mesh (500 and 1000 lpi) with respect to the strip

### 4.9.3.3. About the experimental result

The simulation has demonstrated that the error on the difference $y_{\text{meas}2} - y_{\text{meas}1}$ is equal to the quadratic sum of 15.6 $\mu m$ ($\sqrt{2} \times 11.0 \mu m$) with a systematic error which depends on the track intercept (which corresponds to any value of $\delta_c$). The experimental result, 17 $\mu m$, corresponds to a misalignment of the mesh and strips of the second detector with respect to the first one. This misalignment can be evaluated (Fig. 4.54) to 5 $\mu m$. As a conclusion and as long as we do not know the relative position of the mesh with respect to the strips, the spatial resolution in a detector is the quadratic sum of 11.0 $\mu m$ with a systematic error running from 0 to 14 $\mu m$, leaving a quadratic mean value of 9.6 $\mu m$. We can summarize by saying that the measured spatial resolution is $14 \pm 3 \mu m$. Note that with a 1000 lpi mesh, Figs. 4.52, 4.53 (b) and 4.54 show that the effect of the mesh is
highly reduced and become negligible compared to the diffusion or strip pitch contributions.

**Figure 4.53.** The simulated average systematic difference of the intercepts of tracks in the two detectors of a MICROMEGAS doublet as a function of the relative position of the mesh and strip in one detector for different relative positions of the mesh and strip in the other one, in case of a: (a) 500 lpi, (b) 1000 lpi mesh.

**Figure 4.54.** The simulated error on the difference of the intercepts of tracks in the two detectors of a MICROMEGAS doublet as a function of the position of the mesh and strip in one detector with respect to the other one, in case of a 500 lpi and 1000 lpi mesh.
4.9.3.4. The simulated dependence of the spatial resolution accuracy with the track angle

By simulation we estimate the accuracy of the spatial resolution with the incident track angle on the detector plane (0 deg at normal incidence). With CF₄ (80%) and isobutane (20%) as gas mixture, 3 mm of drift space, 100 μm of strip pitch, and a 1000 LPI mesh, the simulated accuracy on the spatial resolution grows from 11 to 70 μm when the track angle increases from 0° to 10° of incidence as shown in Fig. 4.55 (a). An improvement can be obtained by reducing the drift space as much as possible without a significant loss of efficiency. Fig. 4.55 (b) indicates that a drift space as small as 0.75 mm could give an efficiency greater than 0.98, whereas with 0.5 mm the efficiency does not exceed 0.96.

One can conclude that for a drift space of 0.75 mm we combine a high efficiency and a limited degradation of the resolution with the track angle, as shown in Fig. 4.55 (a). Compared with the 3 mm case, the resolution is better at large angles but worse at small angles. So the choice of the depth of the drift space is crucial and experiment dependent.

4.9.4. Conclusion

With MICROMEGAS, the best spatial resolution which has been obtained so far for tracks at normal incidence is 14 ± 3 μm. Such a spatial resolution is a world record for a gaseous detector at atmospheric pressure. This result was obtained with a CF₄ + Isobutane (80:20) gas mixture and an electric drift field higher than 1.5 kV/cm. With such a gas, the accuracy was limited by the pitch of the micromesh. Under the same conditions, we expect that the use of a 1000 lpi mesh instead of the 500 lpi would have given a resolution of 11 μm.

Our simulation predicts that further improvements are still possible. With an optimized strip pitch of 75 μm and a gain of 2000, the mean cluster size would be 3 and the resolution 8.5 μm. Even better accuracy could be reached with the same gas mixture under pressure. At 2 (4) bars, with a strip pitch of 50 (35) μm, a drift gap of 3 mm and a 1500 (2000) lpi micromesh, a resolution of 4.5 (2.5) μm at normal incidence is expected.
4.10. Rate capability and comparison with other detectors

Many tests [2] have shown that the rate capability of the gaseous detector strongly depends on the type of the incident particle beam. Two particular cases will be distinguished.

4.10.1. Rate capability with X-ray particles

This environment is relevant for applications in the medical field [20]. It has been observed that gas gain does not saturate with rate, up to fluxes of $10^9 \text{ mm}^2 / \text{s}$. Systematic studies in the laboratory show that the maximum achievable gas gain decreases with the flux. At a flux of $10^7 \text{ mm}^2 / \text{s}$ and X-rays of 8 keV energy, the gas gain is higher than $10^3$, which allows a full detection efficiency.
4.10.2. Rate capability with charged particles

Investigations, using high flux of incident beta particles, low energy protons or high energy muons, have shown that the detector has a similar behavior to the X-rays, i.e., MICROMEGAS can cope with very high rates of these particles. However, undesirable effect has been observed when the incident beam is composed by high energy hadrons: a high discharge rates proportional to the incident hadron flux.

It is believed that large ionization deposits trigger discharges. These deposits are probably released by recoil nuclei produced by charged particles, especially hadrons, traversing the detector. As ionization losses are proportional to $z^2$, the recoil nuclei resulted from elastic or quasi-elastic interactions, with energy in the MeV region, are quite efficient to produce heavy ionization in the gas. A typical example is a nuclear interaction on the Argon nuclei producing fully ionized Argon nuclei having a 1 MeV kinetic energy. The whole energy will be lost within 100 microns producing about $10^5$ electron-ion pairs in the conversion gap. This enormous quantity of charges is again multiplied by the detector gain in the amplification gap, exceeding the Rather limit (a few $10^7$), thus triggering a breakdown. Taking into account the nuclear collision length for hadrons, which is about $10^5$ cm in Argon, the probability to produce such process in the conversion gap is of the order of $10^{-6}$. This spark probability produces a serious limitation when the detector has to deal with very high hadron flux.

In the case of muons the corresponding cross section is several order of magnitudes lower, therefore the probability to induce sparks is negligible. MICROMEGAS has been tested with $\sim 5 \times 10^7$ muons in a small area (a few cm$^2$) without serious loss of its performance.

It is quite important to note that in MICROMEGAS the induced sparks are not propagating in the whole area of the detector, instead they are limited in area to maximum of a few mm$^2$, and the duration of this phenomena doesn’t exceed 100 ns. The
possible consequences the discharges vary from a dead time to a destruction of the electronics or even the detector itself.

4.11. Radiation resistance

Parallel plate detectors exhibit high radiation resistance [2]. The electric field is homogeneous over the whole amplification gap and accumulation of undesirable effects, like polymerization during the avalanche process, has a little effect. The radiation resistance of the detector has been tested in the laboratory using an intense X-ray generator. With a gas mixture of Argon + 6% Isobutane, the gas gain of the detector remained stable up to a total accumulated charge of 18.3 mC / mm$^2$, which corresponds to about 10 years of LHC operation at the full luminosity, and at 40 cm from the interaction point [20].

4.12. New Developments in MICROMEGAS

4.12.1. Two Dimension Readout

In High-Energy experiments, the two-dimension readout is usually not a crucial demand, because the tracking of charged particles can be performed using several X-Y planes. In some medical applications, employing, for example, the scanner technique, one-dimensional readout is usually sufficient. For example, anode strips pointing to the X-ray source can be used.

In some applications, however, the two dimensional readout is mandatory. A straightforward way is to use the anode pads. Any pad size is compatible with MICROMEGAS, but a careful study of the implementation of the electronic chain is required. A drawback is the increase of the number of the read-out channels, especially when a high accuracy is required. In order to decrease the number of electronic channels, a second plane consisting of strips perpendicular to the anode strips is required. Several solutions are presently under investigation. For example, the Nantes group has
investigated a solution based on strips printed on a thin 50 micron FR4 insulator on the backside [24]. These strips are perpendicular to the anode strips that are printed on the upper side of the insulator. The first attempt has shown a quite important suppression of the pick-up signal on the second strip plane. A second study used the resistive anode strips and the charge division method. Recent results obtained in a particle beam are promising. Our group is pursuing a novel way to resolve this problem: micromesh made of strips. The first results are encouraging.

4.12.2. Photo-detector

The first idea is to benefit, as it has been already discussed, from the very high gas gain obtained with He based gas mixtures. Such mixtures have a low sensitivity to the ionizing particles and are ideal for the photon counting.

The second idea is the design of a special configuration of the photo-cathode, deposited on top of the metallic grid of MICROMEGAS, that will allow achieving high gas gains, and at the same time fully suppressing the dangerous effect of photon feedback. The first results are encouraging; the detection efficiency of single electrons is close to 100% and the time resolution in the sub-nanosecond range [25], is below 1 ns.

4.12.3. Micro-TPC

It is a new concept of a TPC-like structure, surrounding the beam pipe, mounted very close (about 1 cm) from the interaction point of a particle collider. The MICROMEGAS detector covers the two end plates of the TPC and has fine anode pads (about 200 micron) as read-out elements.

A sketch of the micro-TPC is shown in Figure 4.56. It illustrates a cylindrical TPC with a 10 mm internal radius, 30 mm external radius and a length of 80 mm.

The total drift length is 40 mm that corresponds to a maximum drift time of 32 ns for CF$_4$ gas mixture or about 100 ns for slower gas fillings (i.e. He mixtures). With a low diffusion gas mixture, crude calculations show that the spatial resolution per each
detected point should be about 20 micron in the end plate and 100 micron in the middle of the TPC. Taking into account the large number of detected clusters (about 100) a mean accuracy of a few microns is expected.

Such performance exceeds by far that of the Silicon detector, where the spatial resolution is limited by the multiple scattering effect occurring in the material of the semiconductor. Moreover the high radiation resistance of MICROMEGAS offers the possibility to install the structure close to the beam pipe (in a more challenging configuration is placed inside the beam pipe), in order to improve the impact parameter resolution.

This is a good challenge for the B-meson tagging and the particle identification. Examples of possible applications are the electron colliders; especially those dedicated to B-physics. The idea can be extended to other accelerators, including the very-high luminosity hadron machines.

Figure 4.56. Schematic of micro-TCP structure readout by a MICROMEGAS micro-pad detector
4.13. Programs for Simulation and Measuring Data

4.13.1. Simulation Programs

The complex structure of electrodes and insulators or dielectrics in MICROMEGAS detectors often makes it impossible to give an appropriate and accurate analytic description of the inner field geometry, which is needed for the drift and multiplication procedure in the simulation. Finite element computations however help to overcome the limitations of analytic approaches at the expense of local accuracy or computation time. A major improvement was done using Monte Carlo techniques to compute the gas transport parameters for electrons. Remarkable agreements with measurements prove that underlying models are adequate to reality. The following software was used to investigate the described subjects:

1. **MAXWELL**: Commercial software package including 2D and 3D electrostatic and magnetostatic field solutions with the help of Finite Element Methods (FEM) [17].
2. **MAGBOLTZ**: Computation of transport properties for electrons in various gases using Monte Carlo techniques [18].
3. **GARFIELD**: Computer, program for the detailed two- and three-dimensional simulation originally of drift chambers, which includes drift of particles, diffusion, avalanching and current induction on signal electrodes [19].

4.13.2. Readout electronics and the Data Acquisition (DAQ)

The Fig. 4.57 & 4.58 show the general layout of the readout and the data acquisition of the MICROMEGAS [21, 23, 24]. The charge on the X or Y strips is read out with the help of four Front End (FE) electronic cards based on the Gasplex chip. Each FE card integrates 96 signals (96 strips) and operates at a maximum clock speed of 1MHz. The cards are controlled by a CAEN sequencer with two CRAMS modules (CAEN Readout for Analog Multiplexed Signals) in a VME crate and are powered by a 6V power supply (positive and negative). The Sequencer provides the proper timing signals (Clock, Track and Hold, Clear or Reset) to the FE cards. The CRAM modules integrate and store the
total charge of each channel indicated by the signal provided by the FE cards until the software reads the data and transfers them to the PC for permanent storage and analysis.

**Figure 4.57.** 2D readout board glued on low intrinsic radiation Plexiglas substrate

**Figure 4.58.** Left: The detector on the laboratory bench, equipped with the four electronic cards and the gas pipes. The strip planes and the transfer lines can be distinguished. Right: (above) Scheme of the X-Y strip plane with the readout transfer lines that end to the four connector pads, (below) The general layout of the Data Acquisition

The signal for triggering the MICROMEGAS device is obtained through the use of a preamplifier (an ORTEC 142B), which provides the high voltage for the micromesh cathode as well. The output of the preamplifier is subsequently shaped and amplified to produce the appropriate trigger signal. The trigger rate of the detector in the experiment is
rather low (1 Hz), hence the zero suppression and pedestal subtraction capabilities of the CAEN modules are not utilized and all strip data are recorded.

The expected signal events (i.e. X-rays) have a characteristic mesh pulse that will be useful to the rejection of unexpected shapes for background events. In this aspect, for this MICROMEGAS detector there is also the recording of the mesh pulse via a high sampling VME Digitizing Board, the MATACQ Board [22]. This board can code 4 analog channels of bandwidth up to 300MHz over 12 bits dynamic range and a sampling frequency reaching up to 2GHz and over 2520 usable points. One of these channels is used to record the time structure of the mesh pulse. The Fig. 4.59 is a schematic of the MICROMEGAS trigger and readout.

Figure 4.59. Schematic view of the Trigger and Readout scheme. The trigger signal comes from the MICROMEGAS mesh. The multiplexed data ($^{55}$Fe source signal -5.9 keV- in this oscilloscope photo) are transferred from the Gasplex card to the VME ADC module where online threshold and pedestal subtraction can be applied -The MATACQ card is noted as ‘FADC’
The data acquisition and monitoring system is based on the LabView software package, of National Instruments, and can run on a PC with either the Linux RedHat 7.3.1 (CERN release) or the Windows 2000 operating system. A dual boot PC is used to connect to the VME Controller and run the data acquisition software. The connection is performed via a PCI-MXI2 card sitting on the PCI bus of the PC, a VME-MXI2 controller card sitting on the VME and a 20m long MXI2 cable connecting these two cards. The DAQ system was chosen to run on Linux because it provides (through CERN) the facilities of the CASTOR automatic data archiving system and the xntp software for the synchronization of the PC clock to the GPS universal time.

The online software is controlled by two LabView virtual instrument modules (VIs) (Fig. 4.60):

- The RunControlAll.vi, which is a state machine module controlling the initialization and start/stop of the run along with the creation of the monitoring processes.
- The RunControlAllMonitor.vi, displays the status of the run and the monitoring processes as well as the values of parameters like run number, event numbers, recording file etc.

The monitoring processes include the Event Display and the Run Monitor. The Event Display displays the X-strip, Y-strip charges and MATACQ digitizing pulse of individual events, with a frequency chosen in advance (Fig. 4.61). The Run Monitor performs an analysis of each event and makes one or two dimensional histograms of X strip or Y strip energies, positions and any other quantities of monitoring interest. All processes run asynchronously and communicate to each other via global variables.

The capability to determine the position with the X-Y readout as well as to detect low energy photons is shown in Fig. 4.62.
Figure 4.60. The modules of the MICROMEAGAS acquisition: the control VIs as described in the text.

Figure 4.61. The on-line view of the data acquisition.
4.13.3. GEANT4 – Monte Carlo Simulation

Geant4 is a toolkit for the simulation of the passage of particles through matter. Its areas of application include high energy, nuclear and accelerator physics, as well as studies in medical and space science. The two main reference papers for Geant4 are published in *Nuclear Instruments and Methods in Physics Research A* 506 (2003) 250-30, and *IEEE Transactions on Nuclear Science* 53 No. 1 (2006) 270-278.

Geant4 is developed and maintained by an international collaboration of physicists and computer scientists. The open and collaborative relationship between the development team and its user communities has led to a two-way transfer of technology, since users in fields other than particle physics actively contribute.

4.13.3.1. Applications

Geant4 has applications in many fields: High Energy Physics, Space and Radiation, Medical and Technology Transfer (Fig. 4.63).
Particle Physics – High Energy Physics: the simulation programs play a fundamental role in optimizing the design of particle physics experiments. For the development of reconstruction programs they provide the necessary input in the form of simulated raw data. In the analysis process they are required to understand systematic effects from detector resolution and acceptance as well as the influence of background processes. (BaBar: The BaBar simulation reproduces in detail the generation of events at the interaction point, the propagation of the resulting particles through the detector and the response of the detector to these particles. Detector response quantities are then used to construct candidate events which may analyzed as if they were real data. ATLAS (LHC): Use of Geant4 in the ATLAS Detector Simulation. OSCAR (LHC): Use of Geant4 in the OSCAR simulation program of the CMS experiment. GAUSS (LHC): Use of Geant4 in the GAUSS simulation program of the LHCb experiment. ALICE (LHC): Use of Geant4 in the ALICE Geant4 Simulation. Fermilab: Use of Geant4 at Fermilab for different applications. ILC: Use of Geant4 for the International Linear Collider project).
- **Space and Radiation Applications:** Simulation is equally important in space-based astroparticle physics. Most space probes have to operate for many years without the possibility of physical repair after launch. It is therefore essential to understand the behaviour of all components in the space environment and in particular the effect of radiation on on-board electronics and detectors. *(European Space Agency: Geant4 Space Users' Home Page, Project Support, XMM-Newton Radiation Environment, Space Environment Information System (SPENVIS), Dose Estimation by Simulation of the ISS Radiation Environment (DESIRE), Simulation of Interactions of Radiation with Biological Systems at the Cellular and DNA Level. QinetiQ: Space Energetic Particle Transport and Interaction Modeling studies (SEPTIMESS), Radiation Effects Analysis Tools (REAT), MUlti-LAyered Shielding SImulation Software (MULASSIS). GLAST: Gamma Ray Large Area Space Telescope).*

- **Medical Applications:** Geant4’s extended set of physics models, handling both electromagnetic and hadronic interactions can be used to address a wide variety of medical applications from conventional photon-beam radiotherapy to brachytherapy (using radioactive sources), and from hadron therapy to boron neutron capture therapy. *(GATE: Geant4 Application for Tomographic Emission).*

- **Technology Transfer**
Chapter 5

MICROMEGAS Application in Medicine

Before we discuss the application of this detector in medicine we discuss its general applications. After that a brief walk through the history of gaseous is mentioned as well as a comparison of the micropattern detectors in the field of Medicine.

5.1. General Applications of MICROMEGAS

The high counting rate, the excellent time resolution and the high accuracy of MICROMEGAS combined with its low cost, radiation resistance and robustness, offer a lot of potential applications in high energy physics as well as in the domain of the X-ray imaging [1]. We now present several applications separated in sections which some of them are according to the kind of particle that is detected.

5.1.1. Charge Particle detection

For the tracking of charged particles in High Energy Physics [2], a spatial resolution as good as possible is needed, often in a high-rate environment for which an occupancy of the detector as low as possible is also needed. The conversion gap appears then to be the main source of degradation of the detector performances.
Several groups, using various MICROMEGAS configurations in terms of strip pitch, amplification gap and operating gas, have investigated the space resolution of the detector, with a conversion gap of 3 mm. The spatial resolution of particles at normal incidence is always limited by the transverse diffusion of the drifting electrons which depends on the gas mixture used. All the results stand under 100 μm (σ). The best accuracy, 12 μm (σ) was obtained using an amplification gap of 100 μm; a strip pitch of 100 μm and a gas mixture of CF₄-isobutane (80–20%) [3]. The conclusion is that the space accuracy of MICROMEGAS can satisfy the needs of most of the high-energy experiments for tracking purposes.

The time jitter of MICROMEGAS depends on several parameters: the fluctuations in the time of arrival of the electrons at the entrance of the holes of the micromesh, the ionization density and the longitudinal diffusion of the electrons. The intrinsic resolution due to the amplification structure is less than 1 ns [4] (see Section 6.1).

There are two different configurations of the detector which are proposed in High Energy Physics experiments: the most common one with a small conversion gap and the TPC mode with a large conversion space. The choice of the gas mixture together with the anode elements and the associated front-end electronics depends on the space or/and time resolution needed and on the particles environment (nature, flux, multiplicity). The use of the amplification in MICROMEGAS allows safe operation in a high rate environment.

In the mode with a small conversion gap, the detector is placed in such a way that most of the charged particles cross the detector plane at small incidence. The MICROMEGAS detector comes in some cases to be competitive with the silicon microstrip detector with quite a few advantages: higher radiation resistance, lower cost and lower detector mass. An example, the Alice experiment at future LHC envisages to implement MICROMEGAS detectors in the front part of the calorimeter. This preshower detector consists of a sandwich of two MICROMEGAS chambers surrounding a passive lead converter [5].
A more advanced project is the tracking of particles in the COMPASS experiment at CERN [6]. In this experiment, the detected particles are mainly muons at high rate \((2 \times 10^7 \, /s)\) with a relatively small amount of hadrons \((5 \times 10^4 \, /s)\). As the limitation in gain with muons is three order of magnitude lower than with hadrons, the panel of gas mixtures, leading safe operation of the detector, is quite broad. Large area MICROMEGAS chambers \(40 \times 40 \, cm^2\) have been constructed and successfully tested in the COMPASS beam [7].

In the mode with a large conversion space, some tests have been made in laboratory using a small TPC chamber coupled to a MICROMEGAS amplification structure, in view to the charged particles detection in the future TESLA accelerator. The use of MICROMEGAS has the advantage to catch the positive ions, created during the avalanche process, in a short time (about 100 \(ns\)) allowing the detection of events with high track multiplicity. In such a detector, the longitudinal diffusion of the gas becomes a crucial limitation of the space resolution. Another great motivation of using MICROMEGAS detector is the reduction of ions escaping into the drift volume. Using large ratio between the amplification and drift field this reduction can be of the order of 99%.

### 5.1.2. Low background applications

Some of the features specific to the MICROMEGAS detectors can be exploited in the search for rare events in neutrino and astroparticle physics [8]. The goal is to detect sporadic signals from a large amount of competing back grounds. In particular, in the CAST experiment at CERN [9], the expected signal comes from solar axions conversion into detectable low-energy photons in the \(keV\) energy range. Background rejection will be essential to achieve the intended sensitivity. A MICROMEGAS detector with anode elements of high granularity can largely reduce the background, exploiting its good space and energy resolution.
5.1.3. Neutrons detection

The use of MICROMEGAS as a neutron profiler, with high transparency and low efficiency, has been demonstrated on a neutron beam [10]. The converter was a thin solid target ($^6$Li or $^{10}$B) deposited on the drift electrode with detection of produced ions and proton recoils in the gas. A high rejection of background gamma rays has been observed. A space resolution better than $400 \mu m$ was obtained with a $3 \mu m$ drift gap and charge preamplifiers. The efficiency is limited by the converters. For the detection of slow neutrons several solutions are under study to reach a full efficiency. A neutron profiler has been installed in the n-TOF [11] neutron beam at CERN for the commissioning of this experiment.

5.1.4. Imaging of photons from visible light to high energy gammas

The main problem is the high-efficiency conversion of the radiations to free electrons. The converters can be solid or gaseous. They have to be compatible with the amplifying gaseous filling and be free of the secondary effects due to photon or ions produced in the electron avalanches. The qualities displayed by MICROMEGAS in position, time and energy resolution would permit various improvements in the detection of electromagnetic radiations. The considerable variability in gas mixtures or pressures opens possibilities to build new instruments for the localization of visible photons or X-rays or high-energy gammas.

5.1.4.1. Detection of visible and VUV light

It is easy in MICROMEGAS, with a proper gas filling like helium and isobutane, to reach amplification factors greater than $10^6$; thus permitting easily the detection of single photoelectrons. A remarkable feature is that the pulse height distribution obtained from the absorption of single VUV photons is a peaked Polya distribution (Fig. 5.1 (a)) as with an high gain multistep avalanche chamber [12], thus leading to a great efficiency in the detection of single electrons. This is combined with a very high electric field at the external surface of the grid, very favorable for the extraction of the photoelectrons from the grid of MICROMEGAS covered with a photocathode. It has been tested with CsI,
sensitive in the UV region and may lead to applications for the detection of the light emitted by scintillation crystals like BaF2.

The time jitter in the detection of single photoelectrons from the CsI is smaller than one nanosecond (Fig. 5.1 (b)). This is quite promising for positron emission tomography. For several years an effort is deployed to make photo cathodes sensitive to visible light and compatible with gaseous detectors [13, 14]. The large gains obtained in the multiplication of single electrons and the peaked distribution of the pulse heights is favorable for this application.

Figure 5.1. (a) Detection of single photoelectrons from a CsI photocathode on the micromesh: pulse height distribution (Polya) and (b) Time resolution with single photoelectrons

5.1.4.2. Detection of X-rays and gamma-rays, from 1 KeV to 10 MeV

Several practical applications of wire chambers, for medical and industrial technology have been found. A group at Novossibirsk has demonstrated [15], by installing more than hundred devices in hospitals, that it is possible, with a linear scanning device based on chambers with wires converging to the X-rays focus, to achieve sizeable reduction of doses for the patient. This is of special interest for pregnant women and children. At Schlumberger, a group developed a scanning device for big containers, which is being
used in several harbours or airports [16]. For both applications there is a strong competition from detectors based on semiconductors or scintillators. Using MICROMEGAS the prospects are changing since a serious gain in accuracy, flexibility, cost and efficiency is within reach and is the subject of active research. With a proper construction of MICROMEGAS a good energy resolution has been obtained, 11% (FWHM) at 5.9 keV and 5.4% (FWHM) at 22 keV (Fig. 5.2) [17]. Such a resolution is close to the best obtained with single wire proportional detectors. It differs by only 25% from the accuracy obtained with gaseous scintillation proportional counters [18], which is ultimately limited by the Fano factor.

This opens new opportunities for the applications in medical or industrial radiology, where counting rates larger than the one permitted by wire chambers as necessary and where the better energy resolution can be helpful. The high pressure operation, such as 10 bar of Argon or Xenon is being developed at Biospace, for radiology applications.

Drift volumes filled with high-Z converters have proven their value for efficient detection of the positron annihilation gamma-rays. The drift space can be filled with tubes or parallelepiped voids, where the electric field is parallel to the surface of the converter and the ionization electrons produced in the gas are drifted to the detection volume of wire chambers [19, 20]. In scanning, operating with MICROMEGAS, the wires are replaced by printed strips and it is then possible to improve the position accuracy by an order of magnitude at least.

5.1.4.3. Two dimensional high efficiency gaseous detector

With wire chambers, in specific fields of applications, gaseous detector have already permitted the construction of 2-D imaging detectors: the spherical drift chambers [21] used for soft diffracted X-ray imaging, filled with high pressure xenon, have been shown to be good for the imaging of 60 keV gammas emitted by a short lived isotope, $^{178}$Ta. This has been investigated for nuclear medicine cameras. A MICROMEGAS detector associated with a large spherical drift space may lead to important applications, as the
structural study of large size molecules, using intense X-rays beams from synchrotron radiation.

Such a detector, provided with a MICROMEGAS amplifying chamber and a spherical entrance window, could become an efficient detector both for crystallography and soft X-ray imaging. The expected accuracy is in the 50 \( \mu m \) range and it is in the hands of the electronics designers to get the ultimate accuracy permitted by physics.

**Figure 5.2.** An exceptionally good energy resolution (5.4\% FWMM at 22 keV) with a chamber of small surface (30 cm\(^2\)) and a flat anode

### 5.2. Medical and Biological Applications

Biological imaging studies, such as protein crystallography, using synchrotron radiation beams, are very demanding in terms of particle flux, accuracy, robustness and fast read-out [1]. These studies are often performed with soft X-rays [22, 23], in the range of about 10keV. At higher X-ray energy, however, one requires either heavier gas mixtures or a
high-pressure operation of the detector. The optimal use of the Xenon gas mixtures could be valuable.

Another application of our detector is in medical radiology where the trend is the digital read-out technology in order to replace the photographic film, with improved sensitivity and a spatial resolution comparable to the film. Beta radiography is employed in medical and biological investigations to image human or animal tissues labeled with beta-emitting radionuclides. One approach has found industrial applications consisting of using a multi-step parallel-plate avalanche chamber coupled to an image intensifier and a CCD to read-out the UV light emitted during the avalanche development in the detector [24].

Fig. 5.3 shows a radiography obtained by scanning with MICROMEGAS operated at 4 bar of Xenon gas mixture [2]. This leads to approaches aimed at overcoming the intrinsic low efficiency of gaseous detectors. In scanning devices, a large depth at MICROMEGAS detector, which can be 40 cm without problem, at the highest possible pressure, with strips pointing towards the source, there is room for a serious improvement over wire chambers.

**Figure 5.3.** Radiograph with Micromegas. Scanning radiography obtained at 4 bar of Xe, with 10 cm absorption length and an X-ray tube of 75 keV
At 5 bar; the absorption length in xenon, for 12 keV photons is only 2 mm: It is thus possible to build with MICROMEGAS a simpler detector than the original spherical chamber. With 20 keV photons, in krypton at 10 bar; the absorption length is 3 mm; still making it possible to keep a good accuracy without the exit spherical grid for the drifting electrons. This would be also an ideal tool for mammography.

5.3. Comparison of Micropattern Detectors in Medicine

5.3.1. Gaseous Detectors; A walk through history

The pioneering work done at the beginning of the twentieth century by Thomson, Rutherford, and Geiger [25] just after the discovery of electromagnetic radiation, focused attention to the development of tools to detect radiation. The Single Wire Proportional Counter (SWPC) was one of the essential tools in the early part of the last century.

It was not until the invention by Charpak in 1968 of the Multi-Wire Proportional Chamber (MWPC) (see Fig. 5.4) [26], that a new era was ushered in this field. The main performance features of the MWPC are a space resolution of few hundred μm, two- and three-dimensional localization of incident radiation, excellent energy resolution, rate capabilities of a few kHz/mm². Two track separations of the order of 2 mm have been measured and very large active areas and volumes have been effected over the last three decades for charged particle tracking in medium and high-energy physics. Several excellent review articles have been written on the subject [27].

Some applications of MWPCs may be cited as crystal diffraction, beta chromatography, and dual energy angiography. A low dose X-ray digital radiography scanner based on the MWPC (see Fig. 5.5) was invented by the BINP Novosibirsk group [28], and is presently being applied routinely to examine patients in hospitals in Russia and in France [29, 30]. Fig. 5.6 taken from [30] shows a film of a congenital hip dislocation (Perthes disease) in a 7-year-old boy with satisfactory visualization of the femoral architecture and bone
texture. Upgrades of this device for improved resolution and dose reduction are still in progress [30, 31].

Figure 5.4. Top left: The configuration of the MWPC chamber; a plane of wires –equally spaced– placed in the middle between two cathode planes. Top right: The signal induced in the closest wire and the neighbouring ones. It will be negative in the former one, while in the neighbouring ones positive. Low: With the implementation of a second plane of wires, placed perpendicular to the first one, the spatial information is increased

5.3.1.1. The Micro Strip Generation

An MSGC [32] consists of a pattern of thin anodes and cathode strips laid on an insulating substrate with a pitch of a few hundred μm as sketched in Fig. 5.7 (a). Delimited by a drift electrode above and appropriate potentials applied, the resulting electric field is as shown in Fig. 5.7 (b). The design itself removes the positive ions from
the vicinity of the avalanches -Fig. 5.8, thereby lending high rate capability to this device almost two orders of magnitude higher ($\sim 10^6 / \text{mm}^2 \text{s}$) than MWPCs ($\sim 10^3 / \text{mm}^2 \text{s}$).

The salient features of its operation are localization accuracies $\sim 30 \mu m$ rms, double track resolution of $400 \mu m$, and good energy resolution. Long-term and magnetic field operations have been demonstrated, and apart from high-energy physics [33], these devices have found application in many varied fields of X-ray spectrometry and digital radiography.

The difficulties associated with operating MSGCs [12] began when they were exposed to highly ionizing particles, usually present in a high luminosity machine backgrounds from low energy $g$’s, proton conversions, and nuclear fragments. The highly ionizing particles deposit almost three orders of magnitude more charge in the detection volume as compared to a minimum ionizing particle. In such a circumstance the avalanche to streamer and avalanche to gliding discharge, transitions are more likely damaging the strips (see Fig. 5.9). Several groups dedicated themselves to examine and understand these defects [35-39].
Figure 5.7. (a) The principle (b) field configuration of the MSGC with typical voltage settings

Figure 5.8. Potential surfaces for a typical MSGC by MAXWELL 3D simulation program
What emerges is that the streamer mode of operation in the case of an MWPC is stable due to the fact that the electric field, shown in Fig. 5.9, in the direction of the propagation of the streamer is increasingly weak (between the anode and cathode). In the case of micro-strip detectors, the anode-cathode distance is very small (~50-100 μm). The electric field at the tip of the streamer and that along the surface being high, the streamer is most likely to be followed by a voltage and ionization density dependent discharge. Charging up of surface defects, long-lived excited states, and overlapping avalanches seem to be the culprit lowering the discharge limits of operation of these devices. With this insight, several novel designs appeared on the horizon, some of which are discussed below.

![Figure 5.9. (Left): Damage in an MSGC and (Right): Field along the surface of an MSGC](image)

5.3.1.2. Microdot and MICROMEGAS, Parallel-Plate Type Detectors

Early attempts to create the first micro-needle structure were done by Spindt et al. as early as 1976. Spindt used them successfully to emit electrons towards the phosphor screen in high vacuum, for the purpose of creation of the flat TV screens. Soon, several people [40], for example Va’vra and later on Oed with his own design of the micro-needle array, have tried to use them as the electron amplifiers in the gas.

However, no observable gas gain was measured at normal pressure due to the extremely fine needles (≪1 μm radius and ~1.2 μm high –Fig. 5.10), and resulting in very small
amplification region. Advances in photolithography and application of silicon foundry techniques heralded a new era in the design and fabrication of “Micropattern Detectors”. The Microdot detector (μDOT), sketched in Fig. 5.10, and operated successfully by Biagi [41], is the ultimate gaseous pixel device with anode dots surrounded by cathode rings. The chamber could reach very high gas gains of ~10^6 without a discharge [42].

A very asymmetric parallel plate chamber, the MICROMEGAS detector invented by Charpak and Giomataris [43], takes advantage of the semi-saturation of the Townsend coefficient at high fields ~ 100 kV/cm in several gas mixtures [44], thus being stable in its operation with minimum ionizing particles. Fig. 5.11 shows electrons drifting from the sensitive volume into the amplification volume and an avalanche in the thin multiplying gap. Figure 5.2 shows the excellent energy resolution of this device. Large area, 40 x 40 cm^2 MICROMEGAS detectors are being made and tested [45] for the COMPASS experiment at CERN; see Fig. 5.12.

![Figure 5.10. Left: The concept of the ‘Spind’ cathode detector. Right: Schematic of a Microdot Chamber](image)

This idea of a microgap parallel-plate detector triggered a chain of other inventions. One of them is the microgap Resistive Plate Chambers (RPCs), which immediately began to be used in practice. There are two main developments in this direction: a “timing RPC” [79] and a high position resolution RPC [80, 81]. “Timing” RPCs are parallel plate detectors with the metallic cathodes and anodes made of medium resistivity (ρ ~ 10^9-10^{11} Ω cm) glass. The gap between the cathode and the anode is typically 100-400 μm [79]. The small gap allows one to achieve a very high time resolution, ~50 ps [79, 80]. The main applications today for this RPC are time of flight detectors for high-energy physic
and PET [80-82]. High position resolution RPCs have a slightly different design: their cathodes are made of low resistivity ($\rho \sim 10^4$-$10^8 \, \Omega \cdot \text{cm}$) materials (Si, GaAs) and the anodes –from medium resistivity glasses with metallic strips of 50 $\mu$m pitch (see Fig. 5.13).

**Figure 5.11.** Electron drift lines from a track by GARFIELD simulation program, on the left a photon of 10 MeV passes the detector (in Ar + iC$_2$H$_{10}$ (90:10)) at temperature of 300 K and pressure of 1 atm, while on the right there is the simulation of electron drift and multiplication in MICROMEGAS

**Figure 5.12.** Large size COMPASS MICROMEGAS prototype
The gap between the anode and the cathode is about 100-400 $\mu m$. This allows one to achieve an excellent position resolution- better than 50 $\mu m$ in digital mode [80]. It is remarkable that such RPCs can operate almost at the same gas gains and counting rates as metallic parallel plate avalanche chambers (PPACs) [80]. But in contrast to the metallic PPAC they are spark protected and thus very reliable in exploitation. This is why these detectors almost immediately after their development began to be used for medical imaging applications [83].

Recently, a microgap detector design combining both ideas (MICROMEGAS and microgap RPCs) has been reported [84]. It has a mesh cathode and a medium resistivity anode placed 100 $\mu m$ below the mesh. The main advantages of this detector are that it is spark- protected and has a traditional drift volume (actually this design is similar to one the described in [85], but with a much smaller gap between the cathode’s mesh and the anode plate).

**Figure 5.13.** A schematic draw of a high position resolution RPC

5.3.1.3. **CAT detectors**

A new kind of detector was invented by Lemonnier et al. during the same era as the detectors discussed so far, called the CAT or “Compteur A Trous” [46]. It consists of a narrow hole micro-machined in an insulator metallized on the surface as the cathode.
The metal at the bottom of the hole constitutes the anode. With appropriate potentials and a drift electrode, this scheme acts as a focusing lens for the drifting electrons left in the wake of ionizing radiation as shown in Fig. 5.14. Figure 5.15 shows the typical energy resolution measured with this kind of a device.

Removing the insulator in between leaves the cathode as a micro-mesh, which when placed with a thin gap above the readout electrode, emulates the CAT operation hence named micro-CAT or μCAT, see Fig. 5.16 [47]. This structure could reach gas gain of several $10^4$. With an ingenious scheme of readout from virtual pixels made by current sharing, offering 20 times finer resolution as compared to the readout cell (Fig. 5.17), the μCAT combined with the Virtual pixels is renamed the VIP [48].

### 5.3.1.4. GEM detectors

A new concept of gas amplification was introduced in 1996 by Sauli, the Gas Electron Multiplier (GEM –Fig. 5.18) [49] manufactured by using standard printed circuit wet-etching techniques schematically shown in Fig. 5.19. Comprising a thin (~50 μm) Kapton foil, double-sided clad with Copper, holes are perforated through (Fig. 5.20 & 5.21). The two surfaces are maintained at a potential gradient, thus providing the
necessary field for electron amplification, as shown in Fig. 5.22, and an avalanche of electrons as in Fig. 5.23.

**Figure 5.16.** Microphotograph of a μCAT mesh

**Figure 5.17.** Virtual Pixel readout scheme

**Figure 5.18.** Chemical etching progress of a GEM

**Figure 5.19.** Chemical etching progress of a GEM

**Figure 5.20.** Several holes (with a diameter of the order of 50 μm) are chemically etched through a metal-insulator-metal thin-foil composite

**Figure 5.21.** Streamlines descending through a differential piece of a GEM
Coupled with a drift electrode above and a readout electrode below, it acts as a highly performing micropattern detector. The essential and advantageous feature of this detector is that amplification and detection are decoupled, and the readout board is at zero potential. Permitting charge transfer to a second amplification device, this opens up the possibility of using a GEM in tandem with an MSGC or a second GEM.

**Figure 5.22.** The electric field of a GEM detector, by GARFIELD simulation program

**Figure 5.23.** 3D Electrostatic modeling and avalanche simulation for GEM using Maxwell and Garfield simulation programs

### 5.3.1.5. **Other MICROPATTERNS detectors**

Following the GEM concept and better understanding of the discharge phenomena, several new ‘micro’-detectors have appeared on the scene: Micro-Wire [50], an extension of the μDOT in the third dimension, Micro-Pin Array (MIPA) [51] (see Fig. 5.24 & 5.25), the Micro-Tube [52], Micro-Well [53], Micro-Trench [54] and Micro-Groove [55].

All these authors have tried to minimize the presence of insulators in between the anode and cathode, which is the culprit for gliding discharges along the surface. Figure 5.26 shows a microphotograph of the Micro-Tube detector, with a field map in Fig. 5.27 [52]. Fabricated using combination of laser micro-machining and nickel electroplating, it
consists of ~ 150 μm diameter cathode and an anode tube, which is machined through the well and plated alongside.

Figure 5.24. Gain with a Micro-Wire Detector

Figure 5.25. Complete anode - cathode microstructure

Figure 5.26. The Microtube: the center pin is metallized on the outside, emulating an anode wire

Figure 5.27. Field across a Microtube

This structure results in an electric field that increases rapidly at the anode, similar to the μDOT. However, there is no insulating material on the direct line of sight from the cathode to the anode. These design features are predicted to lead to higher gas gains, better stability with fewer discharges, and the reduction of charging effects. An investigation of the effects of detector geometry upon Microtube performance shows
similar performance to the $\mu$DOT and $\mu$CAT detectors, it also predicts large gains $\sim 10^4$ [52]. As mentioned above, detailed studies have shown that discharges in the presence of highly ionizing particles appear in all micro-pattern detectors at gains of a few thousands.

It is possible to obtain higher gains with poorly quenched gases, since they permit a lower operating voltage, and have a higher diffusion, thus lowering the charge density and photon feedback probability. Combining the MSGC with a GEM, safe operation has been demonstrated up to gains of few ten thousands; $\sim 200$ of such detectors are operating at HERA-B [56]. The DIRAC [57] experiment at the CERN PS also employs MSGC + GEM detectors, which have permitted to improve the momentum resolution by a factor of two; Fig. 5.28 shows an assembled GEM + MSGC.

Putting two GEMs in tandem offers a robust detector, which has been studied in detail [58], and large size devices are being built for the COMPASS experiment [59], Carrying the concept further, adding a third GEM offers an even more stable operation in the worst hadronic beam environment [60]. At gains of $\sim 10^4$, spark probabilities ($\sim 10^{10}$) have been measured. Figure 5.29 shows the probabilities of discharges in a single, double and triple GEM detector [59]. For large sizes, the GEMs are segmented in order to reduce the capacity, thus limiting the energy in a discharge.

![Figure 5.28. An NSGC & GEM assembly for DIRAC tracker](image)
Figure 5.29. The Microtube: the center pin is metallized on the outside, emulating an anode wire

5.3.2. Potential Applications of the “Micro-Generation”

5.3.2.1. MSGCs for X-Ray Imaging

Conventional film radiography has very good spatial resolution, but limited dynamic range. For radiographic film, the storage and display media are the same. For film storage media, dynamic range means that the film image saturates (additional photons do not cause proportional film darkening). The display contrast is fixed at the time of film exposure. One does not see much difference in visible contrast in different parts of a film image, which can have widely different number of photons/pixels. Whereas in a digital system, the storage medium (computer) does not saturate and has infinite dynamic range. The display media being different from storage can be varied at will, i.e., the available display dynamic can be chosen to cover the number \( N \) photons/ pixel from any \( N \) (min) to \( N \) (max). The image can be further enhanced using photon energy information. This has been made possible by using the MSGCs with Xe-CH\(_4\) at high pressures; an example of an image is shown in Fig. 5.30 [61].
5.3.2.2. TPC Readout

For the TESLA experiment at the future Linear Collider [63], a double or triple GEM configuration is under consideration [64, 65] owing to its fast electron signal, minimal magnetic distortion effects, and suppression of ion feedback by design. Figure 5.31 shows some measurements and simulation of the fractional positive ion feedback in a double GEM [62]. Special hexagonal pads are being developed [64] proving unprecedented resolution of 50-60 μm in a TPC, using charge sharing and induction signals.

Figure 5.30. Images of a snail shell with a MSGC operating with Xe-CH₄ at 4 bar [62]

Figure 5.31. Fraction ion feedback in the TPC drift volume (points are measurements [65], and a solid line is a computation)
5.3.2.3. High Counting Rate X-ray Imaging

In the last decade enormous efforts by various research groups and companies have been made to develop digital radiographic devices. The most attractive among them are so-called “photon-counting devices”, which allow one to reduce the dose during the image taking. The high–position resolution RPC is one of the most promising candidates for this application [80]. As an example, Fig. 5.32 shows the image of a fish obtained with a high position resolution RPC and for comparison, the same image obtained by standard film techniques. It is obvious that the quality of the digital image obtained with the RPC is much higher. As was described earlier this device can operate at rates as high as the parallel-plate detectors with metallic electrodes. However, in contrast to the metallic PPAC it is spark-protected and has a position resolution better than 50 μm in digital mode at a counting rate of $10^5 \text{ Hz/mm}^2$ [80].

5.3.2.4. The MICROMEGAS X-ray Gallery

Operating in pure Xenon at atmospheric pressure, the MICROMEGAS detectors have been developed for X-ray imaging. Figure 5.33 [66] shows an example of a vertebra scanned by MICROMEGAS.

5.3.2.5. Protein Crystallography SAXS (Small Angle X-Ray Scattering)

X-ray diffraction studies using MSGCs have yielded rapid analysis of single crystal structures by using the information of position and time of the incident X-rays: crystal structures of organic molecules can be obtained in a matter of minutes [67]. Fast time resolved X-ray diffraction measurements offer a time variation of the SAXS pattern of a protein solution for example, shown in Fig. 5.34, within a frame time of 10 ms.

A diffraction pattern of a lipid membrane is shown in Fig. 5.35. The membrane was made with a VIP detector at Elettra, a synchrotron source in Trieste, Italy. With complex algorithms, especially made for the readout cell border, and superimposing several shots of images, a high degree of detail may be obtained from diffraction patterns (see Fig. 5.36) [48].
Figure 5.32. X-ray images of the fish obtained with the high position resolution RPC and with a standard film

Figure 5.33. The radiographic Image of a human vertebra (70 mm × 25 mm) taken by MICROMEGAS
Figure 5.34. (a) X-ray diffraction patterns of Cytochrome show different concentrations of contamination molecules, (b) X-ray diffraction intensities of the same, measured more than a minute after mixing the solutions

Figure 5.35. A single intensive spot in crystalloagram of a protein crystal, collagen using VIP
Figure 5.36. Scan of the diffraction pattern of the lipid with VIP; the middle part shows linear and logarithmic profiles of the pattern

5.3.2.6. Digital Mammography UV and Visible Photomultipliers using GEM

The importance of early detection of cancer is obvious; small tumors are usually detected in routine radiographic scanning of the human body. Current radiographic equipment is limited in its detection capability by the limited contrast difference exhibited by malignant and benign tissues under given radiation doses. A combination of an x-ray converter, a MSGC, and a visible photocathode, shows great promise for a detector for digital mammography [68]. The essential features are a large flat area and high resolution. With a photocathode (UV, visible) coupled to a micropattern detector, the sealed gas avalanche photomultipliers are being developed for fast imaging of UV and visible light, as well as flat readout devices for scintillator and scintillating fiber arrays, and as medical imaging. Photocathodes have also been attempted by combining them with the Glass Capillary Plate (GCP) detectors [69], with the advantages of reduced photon-photon feedback and the high level of cleanliness necessary for the manufacture of high efficiency of photocathodes/secondary photon emitters.
To this end, single-photon detection has been actively pursued using a CsI photocathode coupled to three or four GEMs in tandem, and very large gains $\sim 10^5$ have been obtained in pure noble gas (Ar), and $10^6$ with an admixture of few percent CH$_4$. [68]. With a small preamplification in the drift region, combined with high diffusion, fully efficient single-photon detection is predicted and measured [70, 71] as demonstrated in Fig. 5.37 & 5.38. Fig. 5.39 shows an example of a large single-photon signal in Ar-CH$_4$ (95:5) with a three GEM combination [72].

**Figure 5.37.** Transmission of single electrons amplified in a GEM

**Figure 5.38.** Full efficiency of single electrons in a GEM with a transmissive photocathode and preamplification in the drift region
5.3.2.7. Cherenkov Ring Imaging (RICH/DIRC)

A feasibility study aimed at improving the detection of photons emitted by Cherenkov light, for example, for an upgrade of the existing SLD CRID at SLAC was made. Using a cascade of four GEMs and an operating gas as pure ethane, very high gains have been observed [70, 73], as shown in Fig. 5.40.

5.3.2.8. Scintillation Light Imaging

A novel application of micro-pattern technology was developed by integrating a MSGC in a gas proportional scintillation counter (GPSC) [74]. Instead of the usual photomultiplier tube, a reflective CsI photocathode was deposited on the microstrip plate surface of the MSGC that serves as the VUV photosensor for the scintillation light from xenon GPSC. This hybrid detector will be used to measure the Lamb shift in muonic hydrogen by detecting the 1.9-keV x-ray from the 2P-1S de-excitation, in a 5-T magnetic field [75].

With a GEM as amplifier and a CCD camera, images of individual projected alpha tracks are visible using the scintillation properties of Ar and CF₄, as shown in Fig. 5.41 [76]. The spectral distribution of the emitted light is analyzed in terms of the number of photons emitted per electron in the visible and near infrared regions \( (400 < \lambda < 1000 \text{ nm}) \). The maximum number of photons emitted decreases with pressure.
Figure 5.40. Gain in a Quadruple GEM cascade in ethane gas at 1 bar and room temperature

Figure 5.41. (Above): Images of alpha tracks taken using the tracking chamber with Ar-CF$_4$ (95:5).
(Below): Scintillation images of alpha tracks ($^{241}$Am) in Ar-CF$_4$
5.3.2.9. X-ray imaging: Radiology and diagnostics

With a GEM + MSGC combination operating in Xe-CH₄ at 4 atm. X-ray images have been taken as an excellent example of imaging for diagnostics with a micro-pattern detector (see Fig. 5.42) [64]. Specialized two-dimensional readout boards have been manufactured using the GEM technology. Operating at ground potential, these boards, in conjunction with GEMs, have been developed for digital absorption radiography [62]. With a pixel size of 50 μm, the image of a mammal (small bat, width 32 mm) is shown in Fig. 5.43 [65].

5.3.2.10. GEM for Plasma Diagnostics

Exploiting the selective sensitivity and the high rate capability of GEM to soft X-rays, imaging the dynamics of fusion plasmas has been attempted by the Frascati and Pisa groups for the Frascati Tokamak Upgrade (FTU). With a GEM and individual pixel readout, time resolved plasma diagnostics are made giving information about temperature and turbulence effects. Figure 5.44 exemplifies the recognition of a steady state and collapsed plasma [78] by integrating counts over 50 μs in four adjacent channels.

5.3.3. Comparing MWPC, GEM and MICROMEGAS Detector

In Fig. 5.45, 5.46 & 5.47 we can observe the gain as well as the discharge probability as a function of differential potential field for GEM (Single, Double, Triple GEM and GEM + PCB) and MICROMEGAS detector, while in Fig. 5.48 & 5.49 the gain stability is shown, and in Fig. 5.50 the spatial resolution for GEM & MICROMEGAS is presented, too. All these results are summarized in the Table 5.1.

Figure 5.42. (Left): A 3 mm × 3 mm 13 kV X-ray absorption radiography of a fish bone taken at 2 atm, (Right): a 3 mm × 10 mm 50 kV X-ray digital image of a mouse
Figure 5.43. (Left): Radiography of a small bat using GEM and 50 μm × 50 μm 2D-readout, (Right): A schematic view of a double-GEM X-ray-imaging detector equipped with a 2D readout board. Photography of a bat’s claw and its X-ray radiographic image, taken at 8 keV, are shown. The localization resolution is of the order of 0.1 mm

Figure 5.44. (Left): Reconstruction of photoelectrons with a GEM + micropixel readout, (Right): Counts integrated in 50 μs for four adjacent pixels as the Frascati Toakamak [77]
5.4. Conclusions

Multiwire chambers have matured since their introduction over the last few decades, with several applications in particle physics and diagnostics of various kinds. The last decade has seen several novel developments in Micropattern Gaseous Detectors of which some have been summarized in this Chapter. Basic understanding of the discharge mechanisms in these devices has also improved, allowing amelioration of their design. Progress in the manufacture of customized readout boards has evolved, revolutionizing the potential applications of these detectors in radiology, diagnostics, astrophysics, and other fields.

Figure 5.45. Gain as well as Discharge probability on 5 MeV a particles as a function of differential potential for GEM with PCB
Figure 5.46. Gain as well as Discharge probability on 5 $MeV$ particles as a function of differential potential for a single, double and triple GEM detector

Figure 5.47. Gain as well as Discharge probability on 5 $MeV$ particles as a function of differential potential for a single, double and triple GEM detector
Figure 5.48. The effective gain as a function of rate for a double GEM detector with PCB

Figure 5.49. The relative gain as a function of rate for a MWPC, MSGC and MICROMEGAS detector
Figure 5.50. The spatial resolution of a GEM (on the left) in the order of 71.3 μm, and of the MICROMEGAS (on the right) in the order of 43 μm.

Table 5.1. Summarizing Table of comparing MWPC, GEM and MICROMEGAS detector

<table>
<thead>
<tr>
<th></th>
<th>MWPC</th>
<th>GEM</th>
<th>MICROMEGAS</th>
</tr>
</thead>
<tbody>
<tr>
<td>Rate Capability</td>
<td>$10^4$</td>
<td>$5 \times 10^5$</td>
<td>$10^6$</td>
</tr>
<tr>
<td>[Hz/ mm$^2$]</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Gain</td>
<td>High $10^6$</td>
<td>Low $10^3$ (single)</td>
<td>High $&gt; 10^5$</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$&gt; 10^5$ (multi GEM)</td>
<td></td>
</tr>
<tr>
<td>Gain Stability</td>
<td>Drops at $10^4$</td>
<td>Stable over $5 \times 10^5$</td>
<td>Stable over $10^6$</td>
</tr>
<tr>
<td>[Hz/ mm$^2$]</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>2D Readout</td>
<td>Not really</td>
<td>Yes and flexible</td>
<td>Yes, not flexible</td>
</tr>
<tr>
<td>Position Resolution</td>
<td>$&gt; 200$</td>
<td>50</td>
<td>Good</td>
</tr>
<tr>
<td>[μm]</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Time Resolution</td>
<td>$\sim 100 \times 10^3$</td>
<td>&lt; 100</td>
<td>4.5</td>
</tr>
<tr>
<td>[ns]</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Magnetic Field Effect</td>
<td>High</td>
<td>Low</td>
<td>Low</td>
</tr>
<tr>
<td>Cost</td>
<td>Expensive, fragile</td>
<td>Cheap, robust</td>
<td>Cheap, robust</td>
</tr>
</tbody>
</table>
Conclusions

A revolution is currently taking place in the development of gaseous detectors of photons and particles. Parallel plate-type and wire-type detectors, which dominated for years in high energy and space flight experiments, are now being replaced by recently invented micropattern gaseous detectors. Since these detectors are cheap, can operate at relatively high gains and have very good position resolutions, they may compete with other types of detectors, for example with solid state detectors.

MICROMEGAS detector is a gaseous detector, which its main characteristics have been described. The measured performances of MICROMEGAS detector can be summarized by:

- High efficiency for minimum ionizing particles with large plateau has been measured.
- Spatial resolution of 12 μm has been achieved in CF₄ gas.
- A time resolution of 4.5 ns has been measured in a particle beam. Using fast current preamplifiers and CF₄ mixture, a time resolution of 1 ns is expected for minimum ionizing particle detection.
- The detector is able to cope with X-ray fluxes, as high as 10⁷/mm²/s. Similar results are obtained when ionization particles are electrons or muons. With high energy hadron beam, however, a limitation has been observed due to the production of highly ionizing particles inducing local breakdown. The detector is, however, radiation hard and can cope with a very-large amount of such discharges, without significant dead time and damage of the electronics.
A new photo-detector using the MICROMEGAS structure was presented, giving excellent single electron efficiency (close to 100%) and time resolution below 500 ps.

A new micro-vertex suitable to measure the impact parameter with an accuracy of a few microns was proposed. It is called micro-TPC and can surround the beam pipe close to the interaction point.

The detector is suitable for many applications in physics or for high rate imaging device.

MICROMEGAS has a lot of applications: high energy physics, astrophysics, plasma diagnostics, medicine, biology and industry. These were some examples of a large field of possible applications where the use of MICROMEGAS can simplify the construction and improve the performance in terms of accuracy and readout speed. Table 6.1 summarizes the results obtained and the ultimate performance expected.

<table>
<thead>
<tr>
<th></th>
<th>Measured</th>
<th>Ultimate</th>
</tr>
</thead>
<tbody>
<tr>
<td>Spatial Resolution [μm]</td>
<td>12 μm in CF₄</td>
<td>&lt; 10</td>
</tr>
<tr>
<td>Time Resolution [ns]</td>
<td>4.5</td>
<td>&lt; 1</td>
</tr>
<tr>
<td>Energy Resolution at 5.9 keV (FWHM)</td>
<td>13%</td>
<td>13%</td>
</tr>
<tr>
<td>Signal to Noise for M.I.Ps</td>
<td>&gt; 100</td>
<td>&gt; 100</td>
</tr>
<tr>
<td>Radiation Hard [mC/ mm²]</td>
<td>10 – 30 years of LHC</td>
<td>&gt; 30</td>
</tr>
<tr>
<td>Rise Time of the Fast Signal [ns]</td>
<td>&lt; 1</td>
<td>&lt; 1</td>
</tr>
</tbody>
</table>
The MICROMEGAS detector is improved that can be used in Medicine while it has an excellent spatial resolution which is better than the photographic films used in radiology and not be able to be edited. Moreover the low cost of its structure and its capability to cover large area make it more attractive to be used in Medicine.

Finally, this detector is compared with the other micropattern detectors and from summary Table 5.1 we make the conclusion that it has better properties than the others; hence it is without doubt a promised detector with many applications, and especially in Medical field.
Large Hadron Collider (LHC)

The Large Hadron Collider (LHC – Fig. A.1) [1] at CERN\(^9\) is being built in a circular tunnel 27 km in circumference. The tunnel is buried around 50 to 175 m underground. It straddles the Swiss and French borders on the outskirts of Geneva. The LHC is designed to collide two counter rotating beams of protons or heavy ions. Proton-proton collisions are foreseen at an energy of 7 TeV per beam.

The beams move around the LHC ring inside a continuous vacuum guided by magnets. The magnets are superconducting and are cooled by a huge cryogenics system. The cables conduct current without resistance in their superconducting state. The beams will be stored at high energy for hours. During this time collisions take place inside the four main LHC experiments [2] which are the following:

- **ALICE (A Large Ion Collider Experiment – Fig. A.2)** – The ALICE Collaboration is building a dedicated heavy-ion detector to exploit the unique physics potential

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\(^9\) CERN is the European Organization for Nuclear Research, the world's largest particle physics centre. It sits astride the Franco-Swiss border near Geneva. CERN is a laboratory where scientists unite to study the building blocks of matter and the forces that hold them together. CERN exists primarily to provide them with the necessary tools. These are accelerators, which accelerate particles to almost the speed of light and detectors to make the particles visible. Founded in 1954, the laboratory was one of Europe's first joint ventures and includes now 20 Member States.
of nucleus-nucleus interactions at LHC energies. Our aim is to study the physics of strongly interacting matter at extreme energy densities, where the formation of a new phase of matter, the quark-gluon plasma, is expected. The existence of such a phase and its properties are key issues in QCD for the understanding of confinement and of chiral-symmetry restoration. For this purpose, we intend to carry out a comprehensive study of the hadrons, electrons, muons and photons produced in the collision of heavy nuclei. Alice will also study proton-proton collisions both as a comparison with lead-lead collisions in physics areas where Alice is competitive with other LHC experiments [3]

- **ATLAS (A large Toroidal LHC ApparatuS – Fig. A.3)** – ATLAS is a particle physics experiment that will explore the fundamental nature of matter and the basic forces that shape our universe. Starting in mid-2008, the ATLAS detector will search for new discoveries in the head-on collisions of protons of extraordinarily high energy. ATLAS is one of the largest collaborative efforts ever attempted in the physical sciences. There are 1900 physicists (including 400 students) participating from more than 164 universities and laboratories in 35 countries [4]

- **CMS (Compact Muon Solenoid – Fig. A.4)** – It contains subsystems which are designed to measure the energy and momentum of photons, electrons, muons, and other products of the collisions. The innermost layer is a silicon-based tracker. Surrounding it is a scintillating crystal electromagnetic calorimeter, which is itself surrounded with a sampling calorimeter for hadrons. The main goals of the experiment are: to explore physics at the TeV scale, to discover the Higgs boson, to look for evidence of physics beyond the standard model, such as supersymmetry, or extra dimensions, to study aspects of heavy ion collisions [5]

- **LHCb (Large Hadron Collider beauty experiment).**
Figure A.1. The structure of LHC

Figure A.2. The ALICE Experiment
Figure A.3. The ATLAS Experiment

Figure A.4. The CMS Experiment
A.1. LHC - Challenges in Accelerator Physics

A.1.1. High Luminosity
In the LHC the energy available in the collisions between the constituents of the protons (the quarks and gluons) will reach the TeV range, which is about 10 times that of LEP and the Fermilab Tevatron. In order to maintain an equally effective physics program at a higher energy $E$ the luminosity of a collider (a quantity proportional to the number of collisions per second) should increase in proportion to $E^2$. This is because the De Broglie wavelength associated to a particle decreases like $1/E$ and hence the cross section of the particle decreases like $1/E^2$. Whereas in past and present colliders the luminosity culminates around $L = 10^{32}\text{cm}^{-2}\text{s}^{-1}$, in the LHC it will reach $L = 10^{34}\text{cm}^{-2}\text{s}^{-1}$. This will be achieved by filling each of the two rings with 2835 bunches of $10^{11}$ particles each. The resulting large beam current ($I_b = 0.53\text{ A}$) is a particular challenge in a machine made of delicate superconducting magnets operating at cryogenic temperatures.

A.1.2. The Beam-Beam Effect Limits the Bunch Density
When two bunches cross in the center of a physics detector only a tiny fraction of the particles collide head-on to produce the wanted events. All the others are deflected by the strong electromagnetic field of the opposing bunch. These deflections, which are stronger for denser bunches, accumulate turn after turn and may eventually lead to particle loss. This beam-beam effect was studied in previous colliders, where experience showed that one cannot increase the bunch density beyond a certain beam-beam limit to preserve a sufficiently long beam lifetime. In order to reach the desired luminosity the LHC has to operate as close as possible to this limit. Its injectors, the old PS and SPS, are being refurbished to provide exactly the required beam density.

A.1.3. Collective Instabilities Must Be Controlled
While traveling down the 27 km long LHC beam pipe at a speed close to the speed of light, each of the 2835 proton bunches leaves behind an electromagnetic wake-field which perturbs the succeeding bunches. In this way any initial disturbance in the position
or energy of a bunch is transmitted to its companions, and under certain phase conditions their oscillations can be amplified and lead to beam loss. These collective instabilities can be severe in the LHC because of the large beam current needed to provide high luminosity. Their effect is minimized by a careful control of the electromagnetic properties of the elements surrounding the beam. For instance the convolutions of the thousands of bellows which are used to allow the machine to contract during cool down are shielded from the beam by thin fingers equipped with sliding contacts; the inner side of the stainless steel beam pipe is coated with pure copper to reduce its resistance to beam induced wall currents. However these precautions cannot suppress all instabilities, and sophisticated feedback systems as well as non linear lenses are being designed to damp the remaining ones.

A.1.4. Particles Have to Remain Stable for Long Times

The beams will be stored at high energy for about 10 hours. During this time the particles make four hundred million revolutions around the machine, a truly astronomical number. Meanwhile the amplitude of their natural oscillations around the central orbit should not increase significantly, because this would dilute the beams and degrade luminosity. This is difficult to achieve, since, in addition to the beam-beam interaction already mentioned, tiny spurious non linear components of the guiding and focusing magnetic-fields of the machine can render the motion slightly chaotic, so that after a large number of turns the particles may be lost. Studies concerning the onset of chaos have become very popular recently in many scientific domains: in particular astronomers now believe that planets in the solar system would show chaotic behaviour if observed for millions of years! The designers of particle colliders take part in this widespread effort, which has direct implications in their field. In the LHC the destabilizing effects of magnetic imperfections is more pronounced at injection energy, because the imperfections are larger owing to persistent current effects in the superconducting cables, and also because the beams occupy a larger fraction of the coil cross section. We must evaluate the Dynamic Aperture, the fraction of the coil cross section within which particles remain stable for the required time, and make sure that it exceeds the dimension of the injected beam with a sufficient safety margin. For the time being, no theory can predict with sufficient
accuracy the long term behaviour of particles in non linear fields. Instead we use fast computers to track hundreds of particles step by step through the thousands LHC magnets for up to a million turns. Results are used to define tolerances for the quality of the magnets at the design stage and during production.

**A.1.5. Beam Losses should not Quench the Magnets**

Despite all precautions the beam lifetime will not be infinite, in other words a fraction of the particles will diffuse towards the beam pipe wall and be lost. In this event the particle energy is converted into heat in the surrounding material and this can induce a quench of the superconducting magnets, thus interrupting operation for hours. To avoid this, a collimation system will catch the unstable particles before they can reach the beam pipe wall, so as to confine losses in well shielded regions far from any superconducting element. The LHC combines for the first time a large beam current at very high energy with the most sophisticated superconducting technology. As a consequence it needs a much more efficient collimation system than previous machines.

**A.1.6. The LHC Lattice should be Flexible**

A modern accelerator or collider is a huge investment which must remain a useful research tool for a long time, and therefore should be adaptable to emerging needs. For instance the CERN SPS accelerator was first upgraded into a proton antiproton collider, then a heavy ion accelerator, later a lepton injector for LEP and now a high density proton injector for LHC. The technical choices made in the LHC to deliver high performance while minimizing cost could drastically reduce the adaptability of the machine, since most of its elements are closely packed and embedded in a continuous cryostat. This is borne in mind by the designers, who make all efforts to include as much flexibility as possible in the lattice to allow further upgrades and cope with unpredictable demands.

**A.1.7. Synchrotron Radiation is Significant in the LHC**

In electron-positron colliders the particles loose every second through synchrotron radiation an amount of energy much larger than the beam stored energy. This loss must
be continuously compensated by the RF system, and as a consequence this phenomenon limits the attainable energy while providing damping of particle oscillations. These effects are unimportant in the LHC because owing to the larger mass of the particles the energy radiated during the same time is only a tiny fraction of the beam energy. They will become significant in proton machines at much higher energies (around 100 TeV). However in the LHC the power emitted, about 3.7 kW, cannot be neglected as it has to be absorbed by the beam pipe at cryogenic temperature. This affects the installed power of the refrigeration system and is an important cost issue. In addition the synchrotron light impinges on the beam pipe walls as a large number of hard U.V. photons. These release absorbed gas molecules, which then increase the residual gas pressure, and liberate photo-electrons, which are accelerated across the beam pipe by the strong positive electric field of the proton bunches. These photoelectrons add to the cryogenic load and may induce an instability of transverse coupled bunch modes.
CAST Experiment

The CERN Axion Solar Telescope (CAST – Fig. B.1) aims to shed light on a 30 year old riddle of particle physics by detecting axions originating from the 15 million degree plasma in the Sun’s core.

Axions were proposed as an extension to the Standard Model of particle physics to explain one of the left intriguing problems in quantum chromo-dynamics (QCD), the so-called strong CP problem: the theory predicts the existence of a CP violating term in the standard equations, yet Nature has never exhibited this in any experiment. The most striking proof of this is the neutron electric dipole moment which is expected to be ~10 orders of magnitude larger than its measured upper limit. The most elegant way to solve this problem is to introduce an additional global symmetry (Peccei & Quinn). As a result, the CP violating term is eliminated and the strong CP problem solved. In the real world we do not observe the Peccei-Quinn symmetry which implies that it is spontaneously broken at some energy scale. The associated Nambu-Goldstone boson is called axion.

Several axion models exist but they all have something in common: axion is a neutral, very light particle that interacts very weakly with ordinary matter. Owing to their potential abundance in the early Universe, axions are also leading candidates for the invisible dark matter of the Universe.
Cosmological and astrophysical considerations have already given some constraints on the axion properties (mass and couplings to fermions, nucleons and photons). According to these bounds, axions are allowed to live only in a small window in the mass range of $\mu$eV to some tens of meV. It has to be taken into account that these constraints are rather uncertain, and have not stopped researches from looking for axions exploring even "forbidden" regions.

Searches for solar axions began 20 years ago when the US Brookhaven Laboratory first pointed an axion telescope at the Sun - a highly useful source of weakly interacting particles for fundamental research, as the solar neutrino anomaly amply demonstrates. Axions would be produced in the Sun through the scattering of photons from electric charges - the Primakoff effect - and their numbers could equal those of solar neutrinos. The idea behind the Brookhaven experiment, first proposed by Pierre Sikivie, was to put the Primakoff effect to work in reverse, using a magnetic field to catalyse the conversion of solar axions back into X-rays photons of a few keVs.
The Brookhaven telescope was later joined by another in Tokyo, while other experiments continued to search in different ways. Experiments at Lawrence Livermore Laboratory and Kyoto, for example, search for relic axions from the early Universe. CERN’s NOMAD experiment joined the hunt, looking for axion production in a neutrino beam. Searches based on axion Bragg scattering have been performed by the SOLAX and COSME collaborations using single crystals of germanium in underground laboratories, while optical detection techniques are used by Italy’s INFN laboratory experiment, PVLAS.

This list is not complete, but, taken together, earlier experiments have scanned the kinetic energy range from $10^{-11}$ eV up to $10^{11}$ eV, so far without success. CAST, however, could make a difference because of the length and strength of the magnetic field that it has available by using a prototype magnet for CERN’s LHC collider, enhanced by the use of a focusing X-ray telescope.

**B.1. The Experiment Structure**

The conversion efficiency for axions increases as the square of the product of the transverse magnetic field component and its length. This makes a 9 tesla, 10 m LHC prototype dipole (fig. B.2) magnet with two straight beam pipes ideal for the task, giving a conversion efficiency exceeding that of the two earlier telescopes by a factor of almost 100. CAST’s LHC magnet is mounted on a platform with $\pm 8^\circ$ vertical movement, allowing for observation of the Sun for 1.5 h at both sunrise and sunset. The horizontal range of $\pm 40^\circ$ encompasses nearly the full azimuthal movement of the Sun during the year. The time the Sun is not reachable is devoted to background measurements. At both ends of the magnet, three different detectors are searching for X-rays coming from axion conversions in the magnet when it is pointing to the Sun. As X-ray detectors CAST utilizes:

- an X-ray mirror telescope in combination with a CCD camera
- a MICROMEGAS
- a TPC

The X-ray focusing system and MICROMEGAS are looking for sunrise axions, while the TPC is occupying both bores on the other end and is waiting for sunset axions.

Figure B.2. The cross section of the straight-bore, twin-aperture LHC dipole prototype [CERN]

B.2. The Operation of CAST

The operation of the CAST experiment is foreseen to go in two phases:

- **Phase I** *(completed)*: during 2003 and 2004 the experiment operated with vacuum inside the magnet pipes and explored axion mass range up to 0.02 eV. 2003 data have been analyzed. No signal above background was observed,
implying an upper limit to the axion photon coupling \( g_{a\gamma} < 1.16 \times 10^{-10} \text{ GeV}^{-1} \) (Phys. Rev. Lett. 94, 121301 (2005)). The analysis of 2004 data also shows no signal above background, yielding an improved upper limit to the axion photon coupling of \( g_{a\gamma} < 8.8 \times 10^{-11} \text{ GeV}^{-1} \) (Fig. B.3, Journal of Cosmology and Astroparticle Physics 04, 010 (2007)).

- **Phase II (running):** in order to extend CAST sensitivity to higher axion rest masses, magnet pipes are filled with a gas (in 2005-2006 with \(^4\text{He}\) and in 2007 and on with \(^3\text{He}\)). With \(^4\text{He}\), gas pressure was increased from 0 - 13.4 mbar, whereas after the conversion of the experiment to run with \(^3\text{He}\) in the first half of 2007 followed by commissioning, CAST will be data taking from October 2007 to end of 2010, covering pressures between 13-120 mbar which corresponds to an axion mass sensitivity up to 1.16 eV.

![Figure B.3. CAST exclusion plot](image-url)
CAST is making an important step in solar axion searches:

1. the sensitivity of the experiment is comparable with the limit imposed by astrophysical considerations,

2. in Phase II, CAST is able to enter into the region which is especially favoured by axion models, for the first time for a laboratory experiment.
Other Nuclear Processes in Radioactive Sources

C.1. Annihilation Radiation

Another source of high-energy photons is the annihilation of positrons. If a positron source such as $^{22}$Na is enclosed or allowed to irradiate an absorbing material, the positrons will annihilate with the absorber electrons to produce two photons, each with an energy equal to the electron mass: $511 \text{ keV}$. In order to conserve momentum, these two photons are always emitted in opposite directions. The $\gamma$ spectrum from a thick positron source will thus show a peak at $511 \text{ keV}$ (corresponding to the detection of one of the annihilation photons) in addition to the peaks characteristic of transitions in the daughter nucleus. Figure C.1 shows the spectrum observed with a thick $^{22}$Na source.

C.2. Internal Conversion

While the emission of a $\gamma$-ray is usually the most common mode of nuclear de-excitation, transitions may also occur through internal conversion. In this process, the nuclear excitation energy is directly transferred to an atomic electron rather than emitted as a photon. The electron is ejected with a kinetic energy equal to the excitation energy minus its atomic binding energy. Unlike $\beta$-decay, therefore, internal conversion electrons are
monoenergetic having approximately the same energy as the competing $\gamma$’s, i.e., a few hundred $keV$ to a few $MeV$.

While the $K$-shell electrons are the most likely electrons to be ejected, electrons in other orbitals may also receive the excitation energy. Thus, an internal conversion source will exhibit a group of internal conversion lines, their differences in energy being equal to the differences in the binding energies of their respective orbitals.

Internal conversion sources are one of the few nuclear sources of monoenergetic electrons and are thus very useful for calibration purposes. Some internal conversion sources readily found in the laboratory are given in Table C.1.

<table>
<thead>
<tr>
<th>Source</th>
<th>Energy [keV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{207}$Bi</td>
<td>480, 967, 1047</td>
</tr>
<tr>
<td>$^{137}$Cs</td>
<td>627</td>
</tr>
<tr>
<td>$^{113}$Sn</td>
<td>365</td>
</tr>
<tr>
<td>$^{133}$Ba</td>
<td>266, 319</td>
</tr>
</tbody>
</table>

**Figure C.1.** Gamma-ray spectrum of a $^{22}$Na source as observed with a NaI detector. Because of positron annihilation in the detector and the source itself, a peak at 511 $keV$ is observed corresponding to the detection of one of the annihilation photons.
C.3. Auger Electrons

As in internal conversion, an excitation which arises in the electron shell can also be transferred to an atomic electron rather than to a characteristic x-ray. Such a process can occur after a reaction such as electron-capture, for example. The electrons emitted are called Auger electrons and are monoenergetic. Like internal conversion lines they can occur in groups, however, their energies are more typical of atomic processes being not more than a few keV. They are thus very susceptible to self-absorption and are difficult to detect.

C.4. Neutron Sources

While it is possible to artificially produce isotopes which emit neutrons, natural neutron emitters which can be used practically in the lab do not exist. Laboratory neutron sources, instead, are based on either spontaneous fission or nuclear reactions.

C.4.1. Spontaneous Fission

Spontaneous fission can occur in many transuranium elements with the release of neutrons along with the fission fragments. These fragments, as well, can promptly decay emitting $\beta$ and $\gamma$ radiation. If the fission source is enveloped in a sufficiently thick container, however, much of this latter radiation can be absorbed leaving only the more penetrating neutrons.

The most common neutron source of this type is $^{252}$Cf which has a half-life of 265 years. The energy spectrum of the neutrons is continuous up to about 10 $MeV$ and exhibits a Maxwellian shape. Figure C.2 shows this spectrum. The distribution is described very precisely by the form [2]

$$\frac{dN}{dE} = \sqrt{E} \exp\left(-\frac{E}{T}\right)$$

(C.1)
A more convenient method of producing neutrons is with the nuclear reactions \((\alpha, n)\) or \((\gamma, n)\). Reactions of this type occur with many nuclei, however, only those with the highest yield are used. Such sources are generally made by mixing the target material with a suitably strong \(\alpha\) or \(\gamma\) emitter. The most common target material is beryllium. Under bombardment by \(\alpha\)'s, beryllium undergoes a number of reactions which lead to the production of free neutrons:

\[
a + ^9\text{He} \rightarrow ^{13}\text{C}^* \rightarrow \begin{cases} ^{12}\text{C}^* + n \\ 8\text{Be} + a + n \\ 3a + n \end{cases} \tag{C.2}
\]

Here the excited compound nucleus \(^{13}\text{C}^*\) is formed which then decays through a variety of modes depending on the excitation energy. In general, the dominant reaction is the decay to \(^{12}\text{C}\) or to the 4.44 MeV excited state of \(^{12}\text{C}\). With \(^{24}\text{La}\) as an \(\alpha\)-source, a neutron yield of about 70 neutrons per \(10^6\) \(\alpha\)'s [3] is generally obtained. With \(^{242}\text{Cm}\), which emits \(\alpha\)'s at a higher energy, the yield is \(\approx 106\) neutrons/\(10^6\) – \(\alpha\) [3]. Other \((\alpha, n)\) neutron sources include \(^{238}\text{Pu}/\text{Be}, ^{226}\text{Ra}/\text{Be}\) and \(^{227}\text{Ac}/\text{Be}\). Targets such as B, F, and Li are also used although the neutron yields are somewhat lower. The half-life of these sources, of course, depends on the half-life of the \(\alpha\)-emitter.

For incident \(\alpha\)'s of a fixed energy, the energy spectrum of neutrons emitted in these sources should theoretically show mono energetic lines corresponding to the different transitions which are made. In mixed sources, however, there is a smearing of the alpha-particle spectrum due to energy loss, so that a large smearing in neutron energy results. There is also considerable Doppler broadening which can amount to as much as 2 MeV. Figure C.3 shows the energy spectrum of neutrons for several sources of this type.

---

\(^{10}\) A common method for denoting nuclear reactions is \(A (x, y)B\) where \(x\) is the bombarding particle, \(A\) the target nucleus, \(B\) the resulting nucleus and \(y\) the the outgoing particle or particles. Note that the ingoing and outgoing particles are always on the inside of the parentheses. The abbreviated notation \((x, y)\), therefore, indicates any nuclear reaction in which \(x\) is the incident particle and \(y\) the resulting, outgoing particle.
In the case of the photo-reaction ($\gamma$, n), only two target materials are suitable: beryllium and deuterium. The respective reactions are

\[ ^9\text{Be} + \gamma \rightarrow ^8\text{Be} + n \quad (C.3) \]

\[ ^2\text{H} + \gamma \rightarrow ^1\text{H} + n \quad (C.4) \]

These sources have the advantage of emitting neutrons which are more or less monoenergetic since the $\gamma$’s are not slowed down as in the case of a’s. The neutrons are, of course, not strictly mono energetic if one works out the kinematics; however, the spread is generally small. The disadvantage of these sources is that the reaction yield per $\gamma$ is 1 – 2 orders of magnitude lower than that of the a-type sources. As well, the nonreacting gammas are not absorbed as easily as a-particles, so that these sources are also accompanied by a large background of $\gamma$ radiation.

**Figure C.2.** Neutron energy spectrum from $^{252}\text{Cf}$ [2]. The form of the spectrum can be described by a Maxwellian distribution.
C.5. Source Activity Units

The activity or strength of a radioactive sample is defined as the mean number of decay processes it undergoes per unit time. This is an extrinsic quantity which depends on the amount of source material contained in the sample - the larger the sample the greater the total number of decays. Moreover, it should be noted that the activity of a source is not necessarily synonymous with the amount of radiation emitted per unit time by the source, although it is certainly related to it. For example, some nuclear transformations result in an unstable daughter nucleus which also disintegrates. Its radiations would then appear with the radiation from the original decay, but would not be included in the activity. Similarly, some nuclides decay through several competing processes, for example, $\beta^+$ - emission or electron capture, where only a fraction of the decays appears as a particular emitted radiation. The relation between radiation output and activity, in fact, depends on the specific nuclear decay scheme, and only in the case of a unique radiative transition is the activity identical to the radiation output.
Activity has traditionally been measured in units of *Curies* (Ci). Originally defined as the activity of 1 g of pure radium-226, this unit is equivalent to

\[ 1 \text{Curie (Ci)} = 3.7 \cdot 10^{10} \text{ disintegrations / s (dps)} \]  

(C.5)

This is, in fact, a very large unit and one generally works in the laboratory with sources on the order of tens or hundreds of micro Curies (μCi).

Because of its rather awkward definition in terms of dps, the *Becquerel*, defined as

\[ 1 \text{Becquerel (Bq)} = 1 \text{ disintegration / s} \]  

(C.6)

is now recommended instead.

It is important to distinguish units of *activity* from those of *dose* such as the Gray or the Sievert. These latter units essentially measure the effects of radiation received by an object or person, whereas the Curie or Becquerel is concerned with the disintegrations in the source itself.

### C.6. The Radioactive Decay Law

The *radioactive decay law* was first established experimentally near the beginning of the century by Rutherford and Soddy and states that the activity of a radioactive sample decays exponentially in time. In terms of modern quantum mechanics, this can easily be derived by considering the fact that a nuclear decay process is governed by a transition probability per unit time, \( \lambda \) characteristic of the nuclear species. If a nuclide has more than one mode of decay, then \( \lambda \) is the sum of the separate constants for each mode

\[ \lambda = \lambda_1 + \lambda_2 + \ldots \]  

(C.7)

In a sample of \( N \) such nuclei, the *mean* number of nuclei decaying in a time \( dt \) would then be

\[ dN = -\lambda N dt \]  

(C.8)

where \( N \) is the number of nuclei and \( A \) is the decay constant. We have assumed here that \( N \) is large so that it may be considered as continuous. Eq. (C.8) may be considered as the
differential form of the radioactive decay law. Integrating (C.8) then results in the exponential,

$$N(t) = N(0) \exp(-\lambda t) \quad (C.9)$$

where $N(0)$ is the number of the nuclei at $t = 0$. The exponential decrease in activity of a radioactive sample is thus governed by the constant $A$. In practice, it is more habitual to use the inverse of $\lambda$,

$$\tau_m = 1/\lambda \quad (C.10)$$

which is known as the mean lifetime. This is just the time it takes for the sample to decay to $1/e$ of its initial activity. Equally in use is the half-life, $T_{1/2}$ which is defined as the time it takes for the sample to decay to one-half of its original activity. Thus,

$$\frac{1}{2} = \exp(-\lambda T_{1/2}) \quad (C.11)$$

which implies

$$T_{1/2} = \frac{1}{\lambda} \ln 2 = \tau_m \ln 2 \ . \quad (C.12)$$

**C.6.1. Fluctuations in Radioactive Decay**

Consider now the number of decays undergone by a radioactive source in a period of time $\Delta t$ which is short compared to the half-life of the source. The activity of the source may then be considered as constant. If repeated measurements of the number of decays, $n$, in the interval $\Delta t$ are now made, fluctuations will be observed from measurement to measurement. This is due to the statistical nature of the decay process; indeed, from quantum mechanics we know that the exact number of decays at any given time can never be predicted, only the probability of such an event. From the radioactive decay law, it can be shown, in fact [5] that the probability of observing $n$ counts in a period $\Delta t$ is given by a Poisson distribution,

$$P(n, \Delta t) = \frac{m^n}{n!} \exp(-m) \quad (C.13)$$

where $m$ is the average number of counts in the period $\Delta t$. The standard deviation of the distribution is then
\[
\sigma = \sqrt{m}
\]  
(C.14)

as is characteristic of Poisson statistics.

**C.6.2. Radioactive Decay Chain**

A very often encountered situation is a radioactive decay chain in which a nuclide decays to a daughter nucleus which itself disintegrates to another unstable nucleus and so on. In the simple case of a three-nucleus chain, i.e.,

\[ A \rightarrow B \rightarrow C, \]

where C is stable, application of the radioactive decay law gives the equations

\[
\frac{dN_a}{dt} = -\lambda_a N_a,
\]

\[
\frac{dN_b}{dt} = \lambda_a N_a - \lambda_b N_b,
\]

\[
\frac{dN_c}{dt} = \lambda_b N_b,
\]

where \( \lambda_a \) and \( \lambda_b \) are the corresponding decay constants. For longer chains, the equations for the additional nuclides are derived in the same manner. If initially \( N_b(0) = N_c(0) = 0 \), solution of (C.15) results in

\[
N_a(t) = N_a(0) \exp(-\lambda_a t),
\]

\[
N_b(t) = N_a(0) \frac{\lambda_a}{\lambda_b - \lambda_a} \left[ \exp(-\lambda_a t) - \exp(-\lambda_b t) \right],
\]

\[
N_c(t) = N_a(0) \left\{ 1 + \frac{1}{\lambda_b - \lambda_a} \left[ \lambda_a \exp(-\lambda_b t) - \lambda_b \exp(-\lambda_a t) \right] \right\}.
\]

(C.16)

The behaviour in time of the three nuclear species is graphed in Fig. C.4. Note that the activity of B here is not given by \( dN_b/dt \) but \( \lambda_b N_b \). This is because \( dN_b/dt \) now also includes the rate of B created by A. We might note also that \( N_b \) goes through a maximum. By setting the derivative to zero, we find

\[
t_{\text{max}} = \frac{\ln \lambda_b}{\lambda_b - \lambda_a}.
\]

(C.17)

At this point, the activity of B is a maximum equal to
\[
\lambda_b N_b (t_{\text{max}}) = \lambda_a N_a (t_{\text{max}})
\]

as seen from from (C.16). This is known as ideal equilibrium. At any other time, the ratio of the activity of B to A (or the ratio of any daughter to its immediate parent in longer chains) is given by

\[
\frac{\lambda_b N_b}{\lambda_a N_a} = \frac{\lambda_b}{\lambda_b - \lambda_a} \left[ 1 - \exp\left(-\left(\lambda_b - \lambda_a\right) t\right) \right]
\]

(C.19)

Three cases may be distinguished:

1. If \( \lambda_a > \lambda_b \), then the ratio increases with time.
2. If \( \lambda_a < \lambda_b \), then (1.23) becomes almost constant \( > 1 \) at large \( t \). This is known as transient equilibrium.
3. If \( \lambda_a \ll \lambda_b \), then the ratio rapidly levels off to \( = 1 \) and reaches a state of secular equilibrium.

These three cases are illustrated in Fig. C.5. In secular equilibrium, note that the number of daughter nuclei B stays constant relative to A. This means that the rate of disintegration of B is the same as its rate of creation. An example is the \( \beta^- \)-decay of \( ^{90}\text{Sr} \):

\[
^{90}\text{Sr} \xrightarrow{\beta^- \text{ 28y}} ^{90}\text{Y} \xrightarrow{\beta^- \text{ 64.8y}} ^{90}\text{Zn}
\]

(C.20)

where the end-point energies for the two \( \beta^- \)s are 0.546 MeV and 2.27 MeV respectively. Since the number of \( ^{90}\text{Y} \) nuclei is kept constant by regeneration from \( ^{90}\text{Sr} \), we essentially have a \( ^{90}\text{Y} \) source with a half-life of 28 y rather than 65 h!

**C.6.3. Radioisotope Production by Irradiation**

A useful application of our example above is in the production of radioactive isotopes by irradiation of a stable element. In such a case, we have the nuclear reaction

\[
A(x, y)B \rightarrow C,
\]

where the isotope B is produced and decays to C with some constant \( \lambda_b \). If \( \sigma (A \rightarrow B) \) is the reaction cross section, \( F \), the flux of irradiating particles \( x \) and \( N_a \), the number of nuclei A, then the radioactive decay equations result in
\[
\frac{dN_a}{dt} = -F \sigma (A \rightarrow B) N_a = -\lambda_a N_a
\]
\[
\frac{dN_a}{dt} = -\lambda_a N_a + \lambda_a N_a
\]
which is in strict analogy to our example above. The maximum yield of isotope $B$, therefore, is obtained at a time $t_{\text{max}}$ given by (C.17).

**Figure C.4.** Radioactive decay of a three nucleus chain

**Figure C.5.** Ratio of daughter to parent radionuclide activity. Curve (a) shows the condition known as transient equilibrium, while curve (b) illustrates secular equilibrium
Appendix D

Calculations of Energy Loss

D.1. Energy Loss of Heavy Charged Particles by Atomic Collisions

D.1.1. Bohr’s Calculation - The Classical Case

Consider a heavy particle with a charge \(ze\) [1], mass \(M\) and velocity \(v\) passing through some material medium and suppose that there is an atomic electron at some distance \(b\) from the particle trajectory (see Fig. D.1). We assume that the electron is free and initially at rest, and furthermore, that it only moves very slightly during the interaction with the heavy particle so that the electric field acting on the electron may be taken at its initial position. Moreover, after the collision, we assume the incident particle to be essentially undeviated from its original path because of its much larger mass \((M \gg m_e)\).

Let us now try to calculate the energy gained by the electron by finding the momentum impulse it receives from colliding with the heavy particle. Thus

\[
I = \int F \, dt = e \int E_\perp \, dt = e \int \frac{dx}{\frac{dx}{u}} = e \int \frac{dx}{u} \quad (D.1)
\]

where only the component of the electric field \(E_\perp\) perpendicular to the particle trajectory enters because of symmetry. To calculate the integral \(\int E_\perp \, dx\), we use Gauss’ Law over
an infinitely long cylinder centered on the particle trajectory and passing through the position of the electron. Then

\[ \int E \, 2 \pi b \, dx = 4 \pi z e, \quad \int E \, dx = \frac{2ze}{b}, \]  

so that

\[ I = \frac{2ze^2}{bu} \]  

and the energy gained by the electron is

\[ \Delta E(b) = \frac{I^2}{2m_e} = \frac{2ze^4}{m_e u^2 b^2} \]  

If we let \( N_e \) be the density of electrons, then the energy lost to all the electrons located at a distance between \( b \) and \( b + db \) in a thickness \( dx \) is

\[ -dE(b) = \Delta E(b) N_e \, dV = \frac{4\pi z^2 e^4}{m_e u^2} N_e \frac{db}{b} \, dx. \]  

where the volume element \( dV = 2\pi b \, db \, dx \). Continuing in a straight forward manner, one would at this point be tempted to integrate (D.5) from \( b = 0 \) to \( \infty \) to get the total energy loss; however, this is contrary to our original assumptions. For example, collisions at very large \( b \) would not take place over a short period of time, so that our impulse calculation would not be valid. As well, for \( b = 0 \), we see that (D.4) gives an infinite energy transfer, so that (D.4) is not valid at small \( b \) either. Our integration, therefore, must be made over some limits \( b_{\text{min}} \) and \( b_{\text{max}} \) between which (D.4) holds. Thus,

\[ \frac{-dE}{dx} = \frac{4\pi z^2 e^4}{m_e u^2} N_e \ln \frac{b_{\text{max}}}{b_{\text{min}}}. \]  

To estimate values for \( b_{\text{min}} \) and \( b_{\text{max}} \), we must make some physical arguments. Classically, the maximum energy transferable is in a head-on collision where the electron obtains an energy of \( \frac{1}{2}m_e (2u)^2 \). If we take relativity into account, this becomes \( 2\gamma^2 m_e u^2 \), where \( \gamma = \left(1 - \beta^2\right)^{-1/2} \) and \( \beta = u/c \). Using (D.4) then, we find

\[ \frac{2\pi z^2 e^4}{m_e u^2 b_{\text{min}}} = 2\gamma^2 m_e u^2, \quad b_{\text{min}} = \frac{ze^2}{\gamma m_e u^2}. \]  

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For $b_{\text{max}}$, we must recall now that the electrons are not free but bound to atoms with some orbital frequency $\nu$. In order for the electron to absorb energy, then, the perturbation caused by the passing particle must take place in a time short compared to the period $\tau = 1/\nu$ of the bound electron, otherwise, the perturbation is adiabatic and no energy is transferred. This is the principle of adiabatic invariance. For our collisions the typical interaction time is $t = b/u$, which relativistically becomes $t \Rightarrow t/\gamma = b/(\gamma u)$, so that
\[
\frac{b}{\gamma u} \leq \tau = \frac{1}{\nu}. \quad (D.8)
\]
Since there are several bound electron states with different frequencies $\nu$, we have used here a mean frequency, $\nu$, averaged over all bound states. An upper limit for $b$, then, is
\[
b_{\text{max}} = \frac{\gamma u}{\nu}. \quad (D.9)
\]
Substituting into (D.6), we find
\[
-k \frac{dE}{dx} = \frac{4 \pi \epsilon^2 e^4}{m_e u^2} N_e \ln \frac{\gamma^2 m u^3}{\epsilon^2 \nu}. \quad (D.10)
\]
This is essentially Bohr’s classical formula. It gives a reasonable description of the energy loss for very heavy particles such as the a-particle or heavier nuclei. However, for lighter particles, e.g. the proton, the formula breaks down because of quantum effects. It nevertheless contains all the essential features of electronic collision loss by charged particles.

Figure D.1. Collision of a heavy charge particle with an atomic electron
D.1.2. The Bethe-Bloch Formula

The correct quantum-mechanical calculation was first performed by Bethe, Bloch and other authors. In the calculation the energy transfer is parametrized in terms of momentum transfer rather than the impact parameter. This, of course, is more realistic since the momentum transfer is a measurable quantity whereas the impact parameter is not. The formula obtained is then

$$-rac{dE}{dx} = 2\pi N_a r_e^2 m_e c^2 \rho \frac{Z Z^2}{A \beta^2} \left[ \ln \left( \frac{2m_e \gamma^2 u^2 W_{\text{max}}}{l^2} \right) - 2\beta^2 \right]$$

(D.11)

Equation (D.11) is commonly known as the Bethe-Bloch formula and is the basic expression used for energy loss calculations. In practice, however, two corrections are normally added: the density effect correction $\delta$, and the shell correction $C$, so that

$$-rac{dE}{dx} = 2\pi N_a r_e^2 m_e c^2 \rho \frac{Z Z^2}{A \beta^2} \left[ \ln \left( \frac{2m_e \gamma^2 u^2 W_{\text{max}}}{l^2} \right) - 2\beta^2 - \delta - 2\frac{C}{Z} \right]$$

(D.12)

with

$$2\pi N_a r_e^2 m_e c^2 = 0.1535 \text{MeV cm}^2 / \text{g}$$

$r_e$: classical electron radius

$Z$: charge of incident particle in units of $e$

$\rho$: density of absorbing material

$\beta = u / c$ of the incident particle

$N_a$: Avogadro’s number

$= 6.022 \cdot 10^{23} \text{mol}^{-1}$

$\gamma = 1 / \sqrt{1 - \beta^2}$

$\delta$: density correction

$I$: mean excitation potential

$C$: shell correction

$Z$: atomic number of absorbing material

$A$: atomic weight of absorbing material

$W_{\text{max}}$: maximum energy transfer in a single collision.

The maximum energy transfer is that produced by a head-on or knock-on collision. For an incident particle of mass $M$, kinematics gives

$$W_{\text{max}} = \frac{2m_e c^2 \eta^2}{1+2s \sqrt{1+\eta^2 + s^2}},$$

(D.13)

where $s = m_e / M$ and $\eta = \beta \gamma$. Moreover, if $M \gg m_e$, then
\[ W_{\text{max}} = 2 m_e c^2 \eta^2. \]

### D.1.2.1. The Mean Excitation Potential

The mean excitation potential, \( I \), is the main parameter of the Bethe-Bloch formula and is essentially the average orbital frequency \( \bar{\nu} \) from Bohr’s formula times Planck’s constant, \( h \bar{\nu} \). It is theoretically a logarithmic average of \( \nu \) weighted by the so-called oscillator strengths of the atomic levels. In practice, this is a very difficult quantity to calculate since the oscillator strengths are unknown for most materials. Instead, values of \( I \) for several materials have been deduced from actual measurements of \( dE/dx \) and a semi-empirical formula for \( I \) vs. \( Z \) fitted to the points. One such formula is

\[
\frac{I}{Z} = 12 + \frac{7}{Z} \text{eV} \\
\frac{I}{Z} = 9.76 + 58.8Z^{-1.19} \text{eV}
\]

(D.14)

It has been shown, however, that \( I \) actually varies with \( Z \) in a more complicated manner [2]. In particular, there are local irregularities or wiggles due to the closing of certain atomic shells. Improved values of \( I \) are given in Table D.1 for several materials [2, 3].

### D.1.2.2. The Shell and Density Corrections

The quantities \( \delta \) and \( C \) are corrections to the Bethe-Bloch formula which are important at high and low energies respectively.

The density effect arises from the fact that the electric field of the particle also tends to polarize the atoms along its path. Because of this polarization, electrons far from the path of the particle will be shielded from the full electric field intensity. Collisions with these outer lying electrons will therefore contribute less to the total energy loss than predicted by the Bethe-Bloch formula. This effect becomes more important as the particle energy increases, as can be seen from the expression for \( b_{\text{max}} \) in (D.9). Clearly as the velocity increases, the radius of the cylinder over which our integration is performed also increases, so that distant collisions contribute more and more to the total energy loss. Moreover, it is clear that this effect depends on the density of the material (hence the
term "density" effect), since the induced polarization will be greater in condensed materials than in lighter substances such as gases. A comparison of the Bethe-Bloch formula with and without corrections is shown in Fig. D.2.

Values for $\delta$ are given by a formula due to Sternheimer:

$$\delta = \begin{cases} 
0 & X < X_0 \\
4.6052 \, X + C_0 + a \left( X_1 - X \right)^m & X_0 < X < X_1 \\
4.6052 \, X + C_0 & X > X_1, 
\end{cases} \quad (D.15)$$

where $X = \log (\beta \gamma)$.

The quantities $X_0$, $X_1$, $C_0$, $a$ and $m$ depend on the absorbing material. The parameter $C_0$ is defined as

$$C_0 = - \left( 2 \ln \frac{I}{h v_p} + 1 \right) \quad (D.16)$$

where $h v_p$ is the so-called plasma frequency of the material, i.e.,

![Figure D.2](image-url)

Figure D.2. Comparison of the Bethe-Block formula with and without the shell and density corrections. The calculation shown here is for copper.
\[ v_p = \sqrt{\frac{N_e e^2}{\pi m_e}} = \sqrt{80.617 \cdot 10^6 \text{ cm}^2 N_e \text{ Hz}} \]  

(D.17)

where \( N_e \) (density of electrons) = \( N_a \rho Z / A \). The remaining constants are determined by fitting (2.30) to experimental data. Values for several materials are presented in Table D.1 [3].

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<th>( -C_0 )</th>
<th>( a )</th>
<th>( m )</th>
<th>( X_1 )</th>
<th>( X_0 )</th>
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<td>BGO</td>
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<td>3.08</td>
<td>3.78</td>
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<tr>
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<td>0.1610</td>
<td>3.24</td>
<td>2.49</td>
<td>0.1464</td>
</tr>
</tbody>
</table>

The shell correction accounts for effects which arise when the velocity of the incident particle is comparable or smaller than the orbital velocity of the bound electrons. At such energies, the assumption that the electron is stationary with respect to the incident particle is no longer valid and the Bethe-Bloch formula breaks down. The correction is generally
small as can be seen in Fig. D.2. We give here an empirical formula [4] for this correction, valid for $\eta \geq 0.1$:

$$C(I, \eta) = \left( 0.422377 \eta^{-2} + 0.0304043 \eta^{-4} - 0.00038106 \eta^{-6} \right) \cdot 10^{-6} I^2 +$$

$$+ \left( 3.850190 \eta^{-2} - 0.1667989 \eta^{-4} + 0.00157955 \eta^{-6} \right) \cdot 10^{-9} I^3$$

(D.18)

where $n = \beta \gamma$ and $I$ is the mean excitation potential in eV.

**D.1.2.3. Other Corrections**

In addition to the shell and density effects, the validity and accuracy of the Bethe-Bloch formula may be extended by including a number of other corrections pertaining to radiation effects at ultra relativistic velocities, kinematic effects due to the assumption of an infinite mass for the projectile, higher-order QED processes, higher-order terms in the scattering cross-section, corrections for the internal structure of the particle, spin effects and electron capture at very slow velocities. With the exception of electron-capture effects with heavy ions, these are usually negligible to within $\pm 1\%$.[5, 6].

**D.1.3. Energy Dependence**

An example of the energy dependence of $dE/dx$ is shown in Fig. D.3 which plots the Bethe-Bloch formula as a function of kinetic energy for several different particles. At non-relativistic energies, $dE/dx$ is dominated by the overall $1/\beta^2$ factor and decreases with increasing velocity until about $u = 0.96c$, where a minimum is reached. Particles at this point are known as minimum ionizing. Note that the minimum value of $dE/dx$ is almost the same for all particles of the same charge. As the energy increases beyond this point, the term $1/\beta^2$ becomes almost constant and $dE/dx$ rises again due to the logarithmic dependence of (D.12). This relativistic rise is cancelled, however, by the density correction as seen in Fig. D.2.

For energies below the minimum ionizing value, each particle exhibits a $dE/dx$ curve which, in most cases, is distinct from the other particle types. This characteristic is often exploited in particle physics as a means for identifying particles in this energy range.
Not shown in Fig. 2.3, is the very low energy region, where the Bethe-Bloch formula breaks down. At low velocities comparable to the velocity of the orbital electrons of the material, $dE/dx$, in fact, reaches a maximum and then drops sharply again. Here, a number of complicated effects come into play. The most important of these is the tendency of the particle to pick up electrons for part of the time. This lowers the effective charge of the particle and thus its stopping power. Calculating this effective charge can be a difficult problem especially for heavy ions.

From Fig. D.3, it is clear that as a heavy particle slows down in matter, its rate of energy loss will change as its kinetic energy changes. And indeed, more energy per unit length will be deposited towards the end of its path rather than at its beginning. This effect is seen in Fig. D.4 which shows the amount of ionization created by a heavy particle as a function of its position along its slowing-down path. This is known as a Bragg curve, and, as can be seen, most of the energy is deposited near the end of the trajectory. At the very end, however, it begins to pick up electrons and the $dE/dx$ drops. This behavior is particularly used in medical applications of radiation where it is desired to deliver a high dose of radiation to deeply embedded malignant growths with a minimum of destruction to the overlaying tissue.

**Figure D.3.** The stopping power $dE/dx$ as function of energy for different particles

**Figure D.4.** A typical Bragg curve showing the variation of $dE/dx$ as a function of the penetration depth of the particle matter. The particle is more ionizing towards the end of its path
D.1.4. Scaling Laws for $dE / dx$

For particles in the same material medium, the Bethe-Bloch formula can be seen to be of the form

$$-\frac{dE}{dx} = z^2 f(\beta), \quad (D.19)$$

where $f(\beta)$ is a function of the particle velocity only. Thus, the energy loss in any given material is dependent only on the charge and velocity of the particle. Since the kinetic energy $T = (\gamma - 1)M c^2$, the velocity is a function of $T / M$, so that $\beta = g(T / M)$. We can therefore transform (D.19) to

$$-\frac{dE}{dx} = z^2 f \left( \frac{T}{M} \right). \quad (D.20)$$

This immediately suggests a scaling law: if we know the $dE / dx$ for a particle of mass $M_1$ and charge $z_1$, then the energy loss of a particle of mass $M_2$, charge $z_2$ and energy $T_2$ in the same material may be found from the values of particle 1 by scaling the energy of particle 2 to $T = T_2 \left( M_1 / M_2 \right)$ and multiplying by the charge ratio $\left( z_2 / z_1 \right)^2$, i.e.,

$$-\frac{dE_2}{dx}(T_2) = -\frac{z_2^2}{z_1^2} \frac{dE_1}{dx} \left( T_2 \frac{M_1}{M_2} \right). \quad (D.21)$$

D.1.5. Mass Stopping Power

When $dE / dx$ is expressed in units of mass thickness, it is found to vary little over a wide range of materials. Indeed, if we make the dependence on material type more evident in the Bethe-Bloch formula, we find

$$-\frac{dE}{d\varepsilon} = -\frac{1}{\rho} \frac{dE}{dx} = \frac{z^2 Z}{A} f(\beta, I), \quad (D.22)$$

where $d\varepsilon = P \, dx$. For not too different $Z$, the ratio $(Z / A)$, in fact, varies little. This is also true of the dependence on $I(Z)$ since it appears in a logarithm, $dE / d\varepsilon$, therefore, is almost independent of material type. A 10 MeV proton, for example, will lose about the same amount of energy in 1 g/cm$^2$ of copper as it will in 1 g/cm$^2$ of aluminium or iron,
etc. As will also be seen, these units are also more convenient when \( dE/dx \)'s are combined for mixed materials.

**D.1.6. \( dE/dx \) for Mixtures and Compounds**

The \( dE/dx \) formula which we have given so far applies to pure elements. What about \( dE/dx \) for compounds and mixtures? Here, if accurate values are desired, one must usually resort to direct measurements; however, a good approximate value can be found in most cases by averaging \( dE/dx \) over each element in the compound weighted by the fraction of electrons belonging to each element (Bragg’s Rule). Thus

\[
\frac{1}{\rho} \frac{dE}{dx} = w_1 \left( \frac{dE}{dx}_1 \right) + \frac{w_2}{\rho_2} \left( \frac{dE}{dx}_2 \right) + \ldots ,
\]

(D.23)

where \( w_1, w_2, \) etc. are the fractions by weight of elements 1, 2, ... in the compound. More explicitly, if \( a_i \) is the number of atoms of the \( i^{th} \) element in the molecule \( M \), then

\[
w_i = \frac{a_i A_i}{A_m}
\]

(D.24)

where \( A_i \) is the atomic weight of \( i^{th} \) element, \( A_m = \sum a_i A_i \).

By expanding (D.23) explicitly and regrouping terms, we can define effective values for \( Z, A, I, \) etc. which may be used directly in (D.12),

\[
Z_{eff} = \sum a_i Z_i , \quad A_{eff} = \sum a_i A_i \quad \text{(D.25)}
\]

\[
\ln I_{eff} = \sum \frac{a_i Z_i \ln I_i}{Z_{eff}} \quad \text{(D.27)}
\]

\[
\delta_{eff} = \sum \frac{a_i Z_i \delta_i}{Z_{eff}} \quad \text{(D.28)}
\]

\[
C_{eff} = \sum a_i C_i \quad \text{(D.29)}
\]

Note here the convenience of working with the mass stopping power, \( 1/\rho (dE/dx) \), rather than the linear stopping power \( dE/dx \).
D.1.7. Limitations of the Bethe-Bloch Formula and Other Effects

The Bethe-Bloch formula as given in (D.12) with the shell and density effect corrections is the usual expression employed in most $dE/dx$ calculations. For elementary particles and nuclei up to the $\alpha$-particle, this formula generally gives results accurate to within a few percent for velocities ranging from the relativistic region down to $\beta = 0.1$. This accuracy may be increased and extended to higher-$Z$ nuclei up to $Z = 26$ by including the charge-dependent corrections mentioned earlier [5, 6].

For $\beta \leq 0.05$, many of the assumptions inherent in the Bethe-Bloch formula are no longer valid even with the corrections. Between $0.01 < \beta < 0.05$, in fact, there is still no satisfactory theory for protons. For heavier nuclei, this is even more the case because of electron capture effects. Some empirical formulae for this energy range may be found in [7]. Below $\beta = 0.01$, however, a successful explanation of energy loss is given by the theory of Lindhard [8].

D.1.8. Channeling

An important exception to the applicability of the Bethe-Bloch formula is in the case of channeling in materials having a spatially symmetric atomic structure, i.e., crystals. This is an effect which occurs only when the particle is incident at angles less than some critical angle with respect to a symmetry axis of the crystal. As it passes through the crystal planes, the particle, in fact, suffers a series of correlated small-angle scatterings which guide it down an open crystal channel. Figure D.5 illustrates this schematically. As can be seen, the correlated scatterings cause the particle to follow a slowly oscillating trajectory which keeps it within the open channel over relatively long distances. The wavelength of the trajectory is generally many lattice lengths long. The net effect of this, of course, is that the particle encounters less electrons than it normally would in an amorphous material (which is assumed by the Bethe-Bloch calculation). When the particle undergoes channeling, therefore, its rate of energy loss is greatly reduced. When working with crystalline materials, it is important therefore to be aware of the crystal
orientation with respect to the incident particles so as to avoid (or achieve, if that is the case) channeling effects.

In general, the critical angle necessary for channeling is small \( (\approx 1^\circ \text{ for } \beta = 0.1) \) and decreases with energy. It can be estimated by the formula [5]

\[
\phi_c = \frac{\sqrt{zZe_0Ad}}{1670}\beta \sqrt{\gamma}
\]  

(D.30)

where \( a_0 \) is the Bohr radius, and \( d \) the interatomic spacing. For \( \phi > \phi_c \), channeling does not occur and the material may be treated as amorphous [8].

![Diagram of channeling in crystalline materials](image)

**Figure D.5.** Schematic diagram of channeling in crystalline materials. The particle suffers a series of correlated scatterings which guides it down an open channel of a lattice.

**D.1.9. Range**

Knowing that charged particles lose their energy in matter, a natural question to ask is: How far will the particles penetrate before they lose all of their energy? Moreover, if we assume that the energy loss is continuous, this distance must be a well defined number, the same for all identical particles with the same initial energy in the same type of material. This quantity is called the *range* of the particle, and depends on the type of material, the particle type and its energy.
Experimentally, the range can be determined by passing a beam of particles at the desired energy through different thicknesses of the material in question and measuring the ratio of transmitted to incident particles. A typical curve of this ratio versus absorber thickness, known as a range number-distance curve, is shown in Fig. 2.7. As can be seen, for small thicknesses, all (or practically all) the particles manage to pass through. As the range is approached this ratio drops. The surprising thing, however, is that the ratio does not drop immediately to the background level, as expected of a well defined quantity. Instead the curve slopes down over a certain spread of thicknesses. This result is due to the fact that the energy loss is not in fact continuous, but statistical in nature. Indeed, two identical particles with the same initial energy will not in general suffer the same number of collisions and hence the same energy loss. A measurement with an ensemble of identical particles, therefore, will show a statistical distribution of ranges centered about some mean value. This phenomenon is known as range straggling: in a first approximation, this distribution is Gaussian in form. The mean value of the distribution is known as the mean range and corresponds to the midpoint on the descending slope of Fig. D.6. This is the thickness at which roughly half the particles are absorbed. More commonly, however, what is desired is the thickness at which all the particles are absorbed, in which ease the point at which the curve drops to the background level should be taken. This point is usually found by taking the tangent to the curve at the midpoint and extrapolating to the zero-level. This value is known as the extrapolated or Practical range (see Fig. D.6).

From a theoretical point of view, we might be tempted to calculate the mean range of a particle of a given energy, $T_0$, by integrating the $dE/dx$ formula,

$$S(T_0) = \int_0^{T_0} \left( \frac{dE}{dx} \right)^{-1} dE.$$  \hfill (D.31)

This yields the approximate pathlength travelled. Eq. (D.31) ignores the effect of multiple Coulomb scattering, however, which causes the particle to follow a zigzag path through the absorber. Thus, the range, defined as the straight-line thickness, will generally be smaller than the total zigzag pathlength.
As it turns out, however, the effect of multiple scattering is generally small for heavy charged particles, so that the total path length is, in fact, a relatively good approximation to the straight-line range. In practice, a semi-empirical formula must be used,

\[ R(T_0) = R_0(T_{\text{min}}) + \int_{T_{\text{min}}}^{T_0} \left( \frac{dE}{dx} \right)^{-1} dE, \]  

(D.32)

where \( T_{\text{min}} \) is the minimum energy at which the \( dE/dx \) formula is valid, and \( R_0(T_{\text{min}}) \) is an empirically determined constant which accounts for the remaining low energy behavior of the energy loss. Results accurate to within a few percent can be obtained in this manner\(^{11}\). Figure D.7 shows some typical range-energy curves for different particles calculated by a numerical integration of the Bethe-Bloch formula. From its almost linear form on the log-log scale, one might expect a relation of the type

\[ R \propto E^b. \]  

(D.33)

This can also be seen from the stopping power formula, which at not too high energies, is dominated by the \( \beta^{-2} \) term,

\[ -dE/dx \propto \beta^{-2} \propto T^{-1}, \]  

(D.34)

where \( T \) is the kinetic energy. Integrating, we thus find

\[ R \propto T^2, \]  

(D.35)

which is consistent with our rough guess. A more accurate fit in this energy range, in fact, gives

\[ R \propto T^{1.75} \]  

(D.36)

which is not too far from our simple calculation. This is only one of many theoretical and semi-empirical formulas which cover many energy ranges and materials [9].

\(^{11}\) We might emphasize here that the range as calculated by (D.32) only takes into account energy losses due to atomic collisions and is valid only as long as atomic collisions remain the principal means of energy loss. At very high energies, where the range becomes larger than the mean free path for a nuclear interaction or for bremsstrahlung emission, this is no longer true and one must take into account these latter interactions as well.
Range-energy relations of this type are extremely useful as they provide an accurate means of measuring the energy of the particles. This was one of the earliest uses of range measurements. As we will see later, they are also necessary for deciding the sizes of detectors to be used in an experiment or in determining the thickness of radiation shielding, among other things.

Because of the scaling of \( dE/dx \), a scaling law for ranges may also be derived. Using (D.21) it is easy then to see

\[
R_2(T_2) = \frac{M_2}{M_1} \frac{Z_2^2}{Z_1^2} R_1 \left( T_2 \frac{M_1}{M_2} \right)
\]  

for different particles in the same medium.

For the same particle in different materials, a rough relation known as the Bragg-Kleeman rule also exists

\[
\frac{R_2}{R_1} = \frac{\rho_2 \sqrt{A_2}}{\rho_1 \sqrt{A_1}},
\]

where \( \rho \) and \( A \) are the densities and atomic numbers of the materials. For compounds, a rough approximation for the range can also be found from the formula

\[
R_{\text{comp}} = \frac{A_{\text{comp}}}{\sum a_i A_i},
\]

where \( A_{\text{comp}} \) is the molecular weight of the compound, \( A_i \) and \( R_i \) are the atomic weight and range of the \( i^{th} \) constituent element, respectively, and \( a_i \) is the number of atoms of the \( i^{th} \) element in the compound molecule.
Figure D.6. Typical range number-distance curve. The distribution of ranges is approximately Gaussian in form.

Figure D.7. Calculated range curves of different heavy particles in aluminium.
D.2. Energy Loss of Electrons and Positrons

D.2.1. Collision Loss

While the basic mechanism of collision loss outlined for heavy charged particles is also valid for electrons and positrons, the Bethe-Bloch formula must be modified somewhat for two reasons. One, as we have already mentioned, is their small mass. The assumption that the incident particle remains undeflected during the collision process is therefore invalid. The second is that for electrons the collisions are between identical particles, so that the calculation must take into account their indistinguishability. These considerations change a number of terms in the formula, in particular, the maximum allowable energy transfer becomes \( W_{\text{max}} = T_e / 2 \) where \( T_e \) is the kinetic energy of the incident electron or positron. If one redoes the calculation, the Bethe-Bloch formula then becomes

\[
-dE/dx = 2\pi N_a \rho \frac{Z}{A} \frac{1}{\beta^2} \frac{2}{(i/m_e c^2)^3} \left[ \ln \frac{\tau^2 (\tau + 2)}{2(i/m_e c^2)} + F(\tau) - \delta - 2 \frac{C}{Z} \right],
\]

where \( \tau \) is the kinetic energy of particle in units of \( m_e c^2 \),

\[
F(\tau) = 1 - \beta^2 + \frac{\tau^2}{8} \frac{(2r + 1) \ln 2}{(\tau + 1)^3}
\]

for \( e^- \)

\[
F(\tau) = 2 \ln 2 - \frac{\beta^2}{12} \left[ \frac{23 + 14}{\tau + 1} + \frac{10}{(\tau + 1)^2} + \frac{4}{(\tau + 1)^3} \right]
\]

for \( e^+ \).

The remaining quantities are as described previously (see D.12–18).

D.2.2. Energy Loss by Radiation: Bremsstrahlung

At energies below a few hundred \( GeV \), electrons and positrons are the only particles in which radiation contributes substantially to the energy loss of the particle. This can easily be seen from the bremsstrahlung cross-sections which we will present in the following section. The emission probability, in fact, varies as the inverse square of the particle...
mass, i.e., $\sigma \propto r_e^2 = \left(\frac{e^2}{mc^2}\right)^2$. Radiation loss by muons ($m = 106\,\text{MeV}$), the next lightest particle, for example, is thus some 40000 times smaller than that for electrons!

Since bremsstrahlung emission depends on the strength of the electric field felt by the electron, the amount of screening from the atomic electrons surrounding the nucleus plays an important role. The cross section is thus dependent not only on the incident electron energy but also on its impact parameter and the atomic number, $Z$, of the material.

The effect of screening can be parametrized by the quantity

$$\xi = \frac{100 m_e c^2 h v}{E_0 E Z^{1/3}} \quad \text{(D.38)}$$

with $E_0$: initial total energy of electron (or positron); $E$: final total energy of electron; $hv$: energy of photon emitted, $(E_0 - E)$. This parameter is related to the radius of the Thomas-Fermi atom and is small, $\xi = 0$, for complete screening and large, $\xi \gg 1$, for no screening.

For relativistic energies greater than a few MeV, the bremsstrahlung cross section is given [10] by the formula

$$d\sigma = 4 Z r_e^2 \frac{d\nu}{\nu} \left\{ (1 + \varepsilon^2) \left[ \frac{\phi_1(\xi)}{4} - \frac{1}{3} \ln Z - f(Z) \right] - \frac{2}{3} \varepsilon \left[ \frac{\phi_2(\xi)}{4} - \frac{1}{3} \ln Z - f(Z) \right] \right\} \quad \text{(D.39)}$$

with $\varepsilon = E_0 / E$, $a = 1/137$, $f(Z)$: Coulomb correction, $\phi_1(\xi)$, $\phi_2(\xi)$ are screening functions depending on $\xi$. This expression is the result of a Born approximation calculation and is not valid at low energies.
For heavy elements \((Z \geq 5)\), the screening functions \(\phi_1\) and \(\phi_2\) are usually calculated using a Thomas-Fermi model of the atom and the values given numerically. A useful approximation accurate to \(\approx 0.5\%\) is given \([11]\) by the empirical formulae
\[
\phi_1(\xi) = 20.863 - 2\ln\left[1 + (0.55846\xi)^2\right] - 4\left[1 - 0.6\exp(-0.9\xi) - 0.4\exp(-1.5\xi)\right] \quad (D.40)
\]
\[
\phi_2(\xi) = \phi_1(\xi) - \frac{2}{3}(1 + 6.5\lambda + 6\xi^2)^{-1},
\]
where
\[
\phi_1(0) = \phi_2(0) + \frac{2}{3} = 4\ln183 \quad \text{as } \xi \to 0
\]
\[
\phi_1(\infty) = \phi_2(\infty) \to 19.19 - 4\ln\xi \quad \text{as } \xi \to \infty.
\]

The function \(f(Z)\) is a small correction to the Born approximation which takes into account the Coulomb interaction of the emitting electron in the electric field of the nucleus. Davies et al. \([12]\) give the formula
\[
f(Z) = a^2\left[1 + a^2\right]^{-1} + 0.20206 - 0.0369a^2 + 0.0083a^4 - 0.002a^6 \quad (D.41)
\]
where \(a = Z/137\).

In the limiting cases of no screening and complete screening, \((D.39)\) can be expressed in simpler analytic forms. For \(\xi \gg 1\) (no screening), \((D.39)\) becomes
\[
d\sigma = 4Zr_e^2a \frac{dv}{v} \left[1 + \frac{e^2 - 2e}{3}\right] \ln \left[\frac{2E_0E}{m_ec^2hv}\right] - \frac{1}{2} - f(Z). \quad (D.42)
\]
For \(\xi = 0\), (complete screening),
\[
d\sigma = 4Zr_e^2a \frac{dv}{v} \left[1 + \frac{e^2 - 2e}{3}\right] \left[\ln \left(183Z^{-1/3}\right) - f(Z)\right] + \frac{e}{9}. \quad (D.43)
\]

The energy loss due to radiation can now be calculated by integrating the cross-section times the photon energy over the allowable energy range, i.e.,
\[
-\left(\frac{dE}{dx}\right)_{\text{rad}} = N \int_0^{\nu_0} hv \frac{d\sigma}{dv} (E_0, \nu) dv \quad (D.44)
\]
with \( N \): number of atoms/cm\(^3\), \( N = \rho N_a / A \); \( \nu_0 = E_0 / \hbar \).

We can rewrite this as

\[
-\left( \frac{dE}{dx} \right)_{\text{rad}} = N E_0 \Phi_{\text{rad}}, \quad \text{where}
\]

\[
\Phi_{\text{rad}} = \frac{1}{E_0} \int h \nu \frac{d\sigma}{dv} (E_0, \nu) dv.
\]  \hspace{1cm} (D.45)

The motivation behind this is that \( d\sigma / dv \) is approximately proportional to \( \nu^{-3} \); the integral \( \Phi_{\text{rad}} \) is therefore practically independent of \( \nu \) and is a function of the material only.

For \( m_e c^2 \ll E_0 \ll 137 m_e c^2 Z^{-1/3} \), \( \xi \gg 1 \), we have no screening, so that integration yields

\[
\Phi_{\text{rad}} = 4 Z^2 r_e^2 a \left( \ln \frac{2E_0}{m_e c^2} - \frac{1}{3} - f(Z) \right). \]  \hspace{1cm} (D.46)

For \( E_0 \gg 137 m_e c^2 Z^{-1/3} \), \( \xi = 0 \), (complete screening)

\[
\Phi_{\text{rad}} = 4 Z^2 r_e^2 a \left( \ln (183 Z^{-1/3}) + \frac{1}{18} - f(Z) \right). \]  \hspace{1cm} (D.47)

At intermediate values of \( \xi \), (D.44) must be integrated numerically.

It is interesting to compare (D.45) to the ionization loss formula in (D.37) \((\text{see Fig. D.8})\). Whereas the ionization loss varies logarithmically with energy and linearly with \( Z \), the radiation loss increases almost linearly with \( E \) and quadratically with \( Z \). This dependence explains the rapid rise of radiation loss.

Another difference is that unlike the ionization loss which is quasicontinuous along the path of the electron or positron, almost all the radiation energy can be emitted in one or two photons. There are thus large fluctuations observed for a beam of monoenergetic electrons or positrons.
D.2.3. Electron-Electron Bremsstrahlung

The above formulae represent the mean energy loss from radiation in the field of the nucleus. There is, however, also a contribution from bremsstrahlung which arises in the field of the atomic electrons. Formulas for electron-electron bremsstrahlung have been worked out by several authors and it can be shown that the cross sections are essentially given by those above except that $Z^2$ is replaced by $Z$. This contribution can thus be approximately taken into account by simply replacing $Z^2$ by $Z(Z+1)$ in all of the above cross-section formulae.

D.2.4. Critical Energy

As we have seen the energy loss by radiation depends strongly on the absorbing material. For each material, we can define a critical energy, $E_c$, at which the radiation loss equals the collision loss. Thus,

$$\left( \frac{dE}{dx} \right)_{rad} = \left( \frac{dE}{dx} \right)_{coll} \quad \text{for} \quad E = E_c.$$

(D.48)
Above this energy, radiation loss will dominate over collision-ionization losses and vice-versa below $E_c$. An approximate formula for $E_c$ [13] is,

$$E_c = \frac{800\text{MeV}}{Z + 1.2}$$  \hspace{1cm} (D.49)

Table D.2 gives a short list of critical energies for various materials so as to give some feeling for the order of magnitudes.

<table>
<thead>
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<th>Material</th>
<th>Critical energy [MeV]</th>
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<td>Pb</td>
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</table>

**D.2.5. Radiation Length**

A similar quantity known as the *radiation length* of the material is even more frequently used. This parameter is defined as the distance over which the electron energy is reduced by a factor 1/e due to radiation loss only. Indeed, if we rearrange (D.45), we get the differential equation

$$-\frac{dE}{dx} = N \Phi_{\text{rad}} \, dx.$$  \hspace{1cm} (D.50)

Considering the high energy limit where collision loss can be ignored relative to radiation loss, $\Phi_{\text{rad}}$ in (D.47) is independent of $E$, so that

$$E = E_0 \exp\left(\frac{-x}{L_{\text{rad}}}\right),$$  \hspace{1cm} (D.51)
where \( x \) is the distance travelled and \( L_{\text{rad}} = 1/N\Phi_{\text{rad}} \) is the radiation length. Using (D.47), we thus find the formula

\[
\frac{1}{L_{\text{rad}}} = \left[ 4Z(Z+1)\frac{\rho N_a}{A} \right] r_e^2 a \left[ \ln(183Z^{-1/3}) - f(Z) \right],
\]

where we have included the contribution from electron-electron bremsstrahlung and ignored the small constant term. Some values of \( L_{\text{rad}} \) are given in Table D.3 for several materials.

**Table D.3. Radiation length for various absorbers**

<table>
<thead>
<tr>
<th>Material</th>
<th>( \text{gm/cm}^2 )</th>
<th>[cm]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Air</td>
<td>36.20</td>
<td>30050.00</td>
</tr>
<tr>
<td>H(_2)O</td>
<td>36.08</td>
<td>36.10</td>
</tr>
<tr>
<td>NaI</td>
<td>9.49</td>
<td>2.59</td>
</tr>
<tr>
<td>Polystyrene</td>
<td>43.80</td>
<td>42.90</td>
</tr>
<tr>
<td>Pb</td>
<td>6.37</td>
<td>0.56</td>
</tr>
<tr>
<td>Cu</td>
<td>12.86</td>
<td>1.43</td>
</tr>
<tr>
<td>Al</td>
<td>24.01</td>
<td>8.90</td>
</tr>
<tr>
<td>Fe</td>
<td>13.84</td>
<td>1.76</td>
</tr>
<tr>
<td>BGO</td>
<td>7.98</td>
<td>1.12</td>
</tr>
<tr>
<td>BaF(_2)</td>
<td>9.91</td>
<td>2.05</td>
</tr>
<tr>
<td>Scint.</td>
<td>43.80</td>
<td>42.40</td>
</tr>
</tbody>
</table>

A useful approximation [13], convenient for quick calculations, is given by

\[
L_{\text{rad}} = \frac{616.4 \text{ g/cm}^2 A}{Z(Z+1)\ln\left(287/\sqrt{Z}\right)},
\]

where \( Z \) and \( A \) are the atomic number and weight of the material respectively. The values obtained are accurate to within 2.5\% except for helium where the result is about 5\% too low.

The usefulness of the radiation length becomes evident when material thicknesses are measured in these units. Clearly, if \( x \) is expressed in units of \( L_{\text{rad}} \), then (D.45) becomes
\[-\left(\frac{dE}{dt}\right) = E_0\]  \hspace{1cm} (D.54)

where \(t\) is the distance in radiation lengths. Thus, the radiation energy loss when expressed in terms of radiation length is roughly independent of the material type.

For compounds and mixtures, the radiation lengths may be computed by applying Bragg’s rule. Expressing \(L_{\text{rad}}\) in mass thickness units, we then have

\[
\frac{1}{L_{\text{rad}}} = w_1 \left( \frac{1}{L_{\text{rad}}^1} \right) + w_2 \left( \frac{1}{L_{\text{rad}}^2} \right) + \ldots,
\]  \hspace{1cm} (D.55)

Where \(w_1, w_2, \ldots\) are the fractions by weight of each element in the mixture as defined in (D.24).

**D.2.6. Range of Electrons**

Because of the electron’s greater susceptibility to multiple scattering by nuclei, the range of electrons is generally very different from the calculated path length obtained from an integration of the \(\frac{dE}{dx}\) formula. Differences ranging from 20 - 400% depending on the energy and material are often found. In addition, the energy loss by electrons fluctuates much more than for heavy particles. This is due to the much greater energy transfer per collision allowed for electrons and to the emission of bremsstrahlung. In both cases, it is possible for a few single collisions (or photons) to absorb the major part of the electron’s energy. This, of course, results in greater range straggling as illustrated by Fig. D.9 which shows some measured range curves.

As for heavy particles, a number of empirical range-energy relations have been formulated. Figure D.10 presents some typical range-energy curves for electrons in various materials as calculated assuming a continuous slowing-down process [15].

**D.2.7. The Absorption of \(\beta\) Electrons**

Because of their continuous spectrum of energies, the absorption of \(\beta\)-decay electrons exhibits a behavior which is very well approximated by an exponential form. This is illustrated in Fig. D.11 which shows the number-distance curves for different absorbers.
plotted on a semi-logarithmic scale. As can be seen, the curves are almost linear and are easily fit by

\[ I = I_0 \exp(-\mu x). \] (D.56)

---

**Figure D.9.** Range number-distance curves for electrons [14]

**Figure D.10.** Range curves for electrons in several materials as calculated in the continuous slowing down approximation [15]

**Figure D.11.** Absorption curves for beta electrons from $^{185}$W [16]
The constant \( \mu \) is known as the \( \beta \)-absorption coefficient and is found to be directly related to the endpoint energy of the \( \beta \)-decay. One of the earliest uses of this behavior was, in fact, to measure \( \beta \) endpoint energies and the thicknesses of thin foils. It is important to note, however, that exponential absorption is not a general characteristic of \( \beta \)-decay. Indeed, this behavior only holds in the case of simple allowed decays. In more complicated forbidden decays where the shape of the \( \beta \)-Spectrum is different, deviations become apparent.


D.3.1. Thick Absorbers: The Gaussian Limit

For relatively thick absorbers such that the number of collisions is large, the energy loss distribution can easily be shown to be Gaussian in form. This follows directly from the Central Limit Theorem in statistics which states that the sum of \( N \) random variables, all following the same statistical distribution, approaches that of a Gaussian-distributed variable in the limit \( N \to \infty \). If we take our random variable to be \( \delta E \), the energy lost in a single atomic collision, and assume that the energy lost in each collision is such that the velocity of the particle is negligibly altered (so that the velocity-dependent collision cross-section stays constant), then the total energy lost is the sum of many independent \( \delta E \), all commonly distributed. Assuming there are a sufficient number of collisions \( N \), then the total will approach the Gaussian form,

\[
f(x, \Delta) \propto \exp \left( \frac{-(\Delta - \overline{\Delta})}{2\sigma^2} \right)
\]

with \( x \): thickness of absorber; \( \Delta \): energy loss in absorber; \( \overline{\Delta} \): mean energy loss; \( \sigma \): standard deviation.

For nonrelativistic heavy particles the spread \( \sigma_0 \) of this Gaussian was calculated by Bohr to be
where \( N_a \) is Avogadro’s number, \( r_e \) and \( m_e \) the classical electron radius and mass, and \( \rho, Z, A \) are the density, atomic number and atomic weight of the material respectively. This formula can easily be extended to relativistic particles

\[
\sigma_0^2 = 4\pi N_a r_e^2 (m_e c^2)^2 \rho \frac{Z}{A} x = 0.569 \rho \frac{Z}{A} x \ [MeV^2],
\]

(D.58)

D.3.2. Very Thick Absorbers

A critical assumption in the above analysis was that the energy loss was small compared to the initial energy so that the velocity change of the particle could be ignored. For very thick absorbers where a substantial amount of energy is lost, this assumption, of course, breaks down [18 - 19].

D.3.3. Thin Absorbers: The Landau and Vavilov Theories

In contrast to the thick absorber case, the distribution for thin absorbers or gases where the number of collisions \( N \) is too small for the Central Limit Theorem to hold is extremely complicated to calculate. This is because of the possibility of large energy transfers in a single collision. For heavy particles, this \( W_{\text{max}} \) is kinematically limited to the expression given in (D.13), while for electrons, as much as one-half the initial energy can be transferred. In this latter case, there is also the additional possibility of a large “One-shot” energy loss from bremsstrahlung as well. While these events are rare, their possibility adds a long tail to the high energy side of the energy-loss probability distribution thus giving it a skewed, asymmetric form. Figure D.12 illustrates this general shape. Note that the mean energy loss no longer corresponds to the peak but is displaced because of the high energy tail. In contrast, the position of the peak now defines the \textit{most probable} energy loss. These two quantities may be used to parameterize the distribution.
Figure D.12. Typical distribution of energy loss in an absorber. Note that is asymmetric with a long high energy tail.

Basic theoretical calculations of this distribution have been carried out by Landau, Symon and Vavilov; each of these, however, has a somewhat different region of applicability. The distinguishing parameter in all these theories is the ratio

$$\kappa = \frac{\bar{\Delta}}{W_{\text{max}}}$$

that is the ratio between the mean energy loss and the maximum energy transfer allowable in a single collision. The mean energy loss may be calculated from the Bethe-Bloch formula, however, for most purposes it is usually approximated by taking the first multiplicative term only and ignoring the logarithmic term, i.e.,

$$\bar{\Delta} = \xi = 2\pi N_e r_c^2 m_e c^2 \rho \frac{Z}{A} \left( \frac{\beta}{\beta_0} \right)^2 x.$$  \hfill (D.61)

Following the literature, we denote this quantity by $\xi$. The thin absorber region is generally taken to be $\kappa < 10$, although for $\kappa > 1$, the distribution already begins to approach the Gaussian limit (see Fig. D.13). By $\kappa > 10$, there, of course, is only a very negligible difference.
D.3.3.1. Landau’s Theory: $\kappa \leq 0.01$

Landau [20] was the first to calculate the energy loss distribution for the ease of very thin absorbers, that is, $\kappa \leq 0.01$. In this theory, Landau makes the assumptions that:

1. the maximum energy transfer permitted is infinite, $W_{\text{max}} \to \infty$, in essence taking $\kappa \to 0$
2. the individual energy transfers are sufficiently large such that the electrons may be treated as free. Small energy transfers from so-called distant collisions are ignored,
3. the decrease in velocity of the particle is negligible, i.e., the particle maintains a constant velocity.

The distribution is then expressed as

$$f(x, \Delta) \frac{\phi(\lambda)}{\xi},$$

where

$$\phi(\lambda) = \frac{1}{\pi} \int_{0}^{\infty} \exp(-u \ln u - u \lambda) \sin \pi u \, du,$$

$$\lambda = \frac{1}{\xi} \left[ \Delta - \xi (\ln \xi - \ln \varepsilon + 1 - C) \right],$$

$$C = \text{Euler's Const} = 0.577 \ldots,$$

and

$$\ln \varepsilon = \ln \frac{(1 - \beta^2) I^2}{2m c^2 \beta^2 + \beta^2}.$$

The quantity $\varepsilon$ essentially represents the minimum energy transfer allowed by assumption 2. The function $\phi(\lambda)$ is a universal function depending only on the parameter and must be evaluated numerically [21-23].

From an evaluation of $\phi(\lambda)$, the most probable energy loss is found to be

$$\Delta_{mp} = \xi \left[ \ln \left( \frac{\xi}{\varepsilon} \right) + 0.198 - \delta \right],$$

where we have also added on the density effect for completeness.
**D.3.3.2. Symon’s Theory and Vavilov’s Theory: Intermediate $\kappa$.**

The region between small covered by Landau and the Gaussian limit is treated by Symon and by Vavilov. Using the limiting distribution derived by Landau, Symon was able to make a number of ingenious approximations in deriving the energy-loss distributions. His results, unfortunately, are expressed in graphic form, which in today’s world of computers, make them inconvenient to use [24].

Vavilov’s theory, in contrast, is along the line of Landau’s formulation and in fact, generalizes the latter’s calculation by taking into account the correct expressions for maximum allowable energy transfer. The latter two assumptions made by Landau are kept however. His results are somewhat more complicated, but reduce to the Landau distribution in the limit $\kappa \to 0$ and to a Gaussian form in the limit $\kappa \to \infty$ [21, 23, 25].

To give an idea of Vavilov’s results, we show Vavilov’s distributions for various values of $\kappa$ in Fig. D.13. These should be compared to the Landau distribution (denoted by $L$) at $\kappa = 0$, also shown in Fig. D.13. Note also how the distribution already resembles a Gaussian form for $\kappa = 1$. In the Gaussian limit, Vavilov gives the variance as

$$\sigma^2 = \frac{\kappa}{\kappa^2} \left( 1 - \beta^2 \right),$$

(D.64)

which agrees with Bohr’s formula for heavy particles in (D.59).

To see how theory compares with experiment, some measured results are also shown in Fig. D.14.
Figure D.13. Vavilov distributions for various $\kappa$. For comparison, Landau’s distribution (denoted by the $L$) for $\kappa = 0$ is also shown [21].

Figure D.14. Comparison of Vavilov’s and Symon’s theories with experiment [21].
I would like to thank my advisor, Pr. Dr. G. Nikiforidis, whose gave me the opportunity to collaborate with the National Technical University of Athens, hence to research this Master Thesis. I would like also to express my gratitude to my supervisor Pr. Dr. Ev. N. Gazis of National Technical University of Athens, without his motivation and encouragement I would not have considered a postgraduate career in research. Thanks to Pr. Dr. Ev. N. Gazis I had the chance to collaborate once again with European Organization for Nuclear Research (CERN) and with National Center of Scientific Research “Demokritos”. I appreciate his vast knowledge and skills, and also his support in many areas and especially in finding scholarship in order to work at CERN.

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