The Development of Two-Phase Xenon Dark Matter Detectors

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A DISSERTATION
PRESENTED TO THE FACULTY
OF PRINCETON UNIVERSITY
IN CANDIDACY FOR THE DEGREE
OF DOCTOR OF PHILOSOPHY

RECOMMENDED FOR ACCEPTANCE
BY THE DEPARTMENT OF
PHYSICS

Adviser: Tom Shutt

June, 2009
Abstract

The nature of dark matter remains one of the great unsolved mysteries of modern physics. The existence of dark matter has been inferred from its gravitational interactions and is strongly supported on theoretical grounds. A primary candidate for the dark matter is the Weakly Interacting Massive Particle (WIMP), which may be an undiscovered particle from the supersymmetric sector. This dissertation describes the research and development in two-phase liquid xenon dark matter detector technology and the results from the full-scale detector XENON10. Two-phase liquid xenon detectors use position sensitivity and simultaneous measurement of light and charge to remove background electron recoil events. The development of this technology has been rapid – the work in this dissertation began in the summer of 2003 when the potential of this technology had yet to be determined, and in early 2008 the XENON10 collaboration published the then world-best upper limit on the spin-independent WIMP-nucleon cross section. The first measurement of the charge based discrimination performance at low energies was achieved in a prototype in early 2005. This prototype also determined the performance of discrimination via scintillation pulse shape. Although pulse shape discrimination was shown to be far weaker than that from charge yield, the combined use of the two methods demonstrated a discrimination power beyond that achieved by either method alone. Alternative detector technologies were also explored. Electron multiplication on wire grids was demonstrated in a two-phase prototype and its discrimination power potential is shown to be near that of the typical electroluminescence charge-readout technique. This could allow for the removal of some or all of the photomultipliers in the detector, which would greatly reduce radioactive backgrounds. The use
of a wavelength shifter was tested in an attempt to improve light collection and was shown to impede charge collection. The magnitude of fluctuations in recombination in electron recoils, which are an important component in charge resolution, was also measured. The final chapter describes the XENON10 detector, its operation at the Gran Sasso National Laboratory, the analysis of the data, and the results of the WIMP search.
Acknowledgements

First and foremost, I would like to thank my advisor, Professor Tom Shutt for providing such incredible scientific insight, encouragement and support throughout the past six years. I would also like to thank Professor Elena Aprile for being so welcoming during my time at Nevis and the other leaders of XENON10, in particular Rick Gaitskell, Karl Giboni, Laura Baudis, Richard Gaitskell, Jose Lopes, Dan McKinsey, and Uwe Oberlack. My work would also not be possible without the help Alexander Bolozdynya, Pawel Majewski, Masaki Yamashita, Pawel Brusov, Kirk McDonald, and Changguo Lu. I would also like to thank Cristiano Galbiati for reviewing my thesis.

I would also like to express special thanks to my fellow flannel-wearing labmate Eric Dahl for helping me maintain my sanity during the endless hours in lab. I wish him success in Chicago and many long nights listening to Wesley Willis and William Hung in lab.

I would also like to express thanks to my numerous colleagues: Alfredo Ferella, Kaixuan Ni, Guillaume Plante, Maria Elena, Angel Manzur, Louis Kastens, Peter Sorensen, Luiz de Viveiros, Simon Fiorucci, Joerg Orboeck, Ana Conceição, Sean Bryan, Jon Leong, Adam Bradley, Annika Peter, Justin Kinney, Nathan Shaman, Jeremy Mock, Sander Zanderbergen, Mike Usowicz, and Sashank Karri.

I am also grateful to my friends back home – Katrina Pagonis, Nick Pagonis, Loretta Tse, Jia Yuan Li, and Michael Wong – for constantly hassling me to finish this thesis. I would especially like to thank you, my darling, for always believing in me and being there for me whenever I needed you. You are the best (and you know it) and I will never forget your artistic renderings of WIMPs.
Finally, I would like thank my family for all their love and support.
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Chapter 1

Dark Matter

1.1 Introduction

Most of the matter in the universe is dark and its nature is unknown. Revealing the nature of dark matter will have enormous implications for our fundamental understanding of the universe. The existence of dark matter was first inferred through its gravitational effects on normal matter and with the advent of precision cosmology, its contribution to the total mass-energy has been pinpointed. This chapter reviews some of the basic evidence and proposed explanations for dark matter.

1.2 Evidence for Dark Matter

Much of the early evidence for dark matter was from observations of the motion of gravitationally bound bodies at the galaxy and cluster level. Recently, evidence has come from more sophisticated techniques such as measurements of the CMB anisotropies, gravitational lensing by galaxy clusters, primordial elemental abundances, and redshift surveys of large-scale structure. Furthermore, these measurements strongly suggest that much of the matter is non-relativistic (cold) and non-baryonic.
1.2.1 Evidence at the Galactic Level

The rotational curves of spiral galaxies provide one of the clearest indications of dark matter. The velocity is determined by measuring the Doppler redshift of atomic transition lines and the common 21 cm line (HI) from neutral hydrogen clouds. The velocity is expected to drop off based on the luminous mass but has been well observed to level off with distance from the galactic center. Assuming that Newtonian dynamics is correct at long distances, the circular velocity at a radius $r$ is

$$v(r) = \sqrt{\frac{GM(r)}{r}},$$

where $M(r)$ is the total mass contained in $r$. The mass is $M(r) = 4\pi \int \rho(r)r^2 dr$. For $v(r)$ to be constant, we must have $M(r) \propto r$ or, equivalently, $\rho(r) \propto r^{-2}$.

Persic et al probed the general mass structure properties of spirals by using the rotation curves of about 1100 spiral galaxies with brightness over a range of 6 magnitudes and distance out to 2 optical radii. Galaxies of different sizes are compared by scaling to the optical radius $R_{opt}$ which is the radius that contains 83% of the light. It has been shown that all spiral rotation curves can be represented by a “universal rotation curve,” which depends only on the luminosity. Figure 1.1 shows the fit of the universal rotation curve to galaxies divided up into 11 luminosity bins. Here, we see that the rotational curves rise sharply at $R/R_{opt} < 1$ and then level off. This study also found that as the luminosity is decreased, the discrepancy increases between luminous matter and matter based on the rotation curve.

The presence of dark matter can also be seen in elliptical galaxies. Elliptical galaxies generally lack neutral hydrogen which is ubiquitous and extensive in spirals, but instead have hot x-ray emitting (thermal bremsstrahlung) gas in hydrostatic equilibrium that provides a means of measuring the dark matter distribution. Dark matter comprises a significant but not dominate portion of the total mass within the half-light radius $r_e$ and increase in dominance at larger radii. An excellent summary on dark matter in elliptical galaxies is given by [76].
Figure 6. Best two-component fits to the universal rotation curve (dotted line: disc; dashed line: halo). The URC beyond $R_{\text{opt}}$ is built by linear extrapolation according to eq.(6). Notice that the extent of the RCs and the smallness of their rms errors limit the uncertainties on the parameters $\beta$ and $a$ to about 10% and 5%, respectively.

The radius encompassing a mean halo overdensity of $<\delta \rho/\rho> = 200$. To smooth the density field we use a top-hat filter: $<\rho>_{R_{200}} = M_{\text{DM}}(R_{200})/(4\pi R_{200}^3)$, with the halo mass $M_{\text{DM}}(R_{200})$ obtained by integrating the central halo density profile of light. These plots are from [1].

Recalling that the mean mass density of the Universe is $\rho = 3H_0^2/(8\pi G)$, then $R_{200}$ is obtained by solving

$$\left< \frac{\delta \rho_H}{\rho} \right>_{R_{200}} = \frac{2}{H_0^2} \left[ V^2 - V_h^2 \right]_{R_{200}} = 200,$$

with $V(R) = V_{\text{URC}}(R)$ for $R \leq 2 R_{\text{opt}}$, and $V(R) = V_{\text{URC}}(2R_{\text{opt}})$ for $R > 2 R_{\text{opt}}$. The quantity $V^2 - V_h^2 = V_h^2$, appearing in (12), can be obtained either directly from (9) and (11) or by using the relation

$$V_h^2(R_{200}) \simeq \left( \frac{L}{L_z} \right) 0.4 \frac{V^2(R_{\text{opt}}) R_{200}}{R_{\text{opt}}},$$

(see Persic & Salucci 1990b): the two estimates are in very good mutual agreement. [In any case, let us remark that $M_{LM} << M_{200}$, with $M_{200} \equiv M_{\text{DM}}(R_{200})$.] In Fig.7 we plot $V(\frac{R}{R_{200}}; L)$: we realize that, when scaled to the DM reference frame, the
1.2.2 Evidence at the Extra-Galactic Level

Dark matter has been also observed in the space between galaxies of groups ($\lesssim 50$ galaxies) and clusters ($\gtrsim 50$). The dark matter content can be measured by virial methods, x-ray and gravitational lensing. The virial theorem allows the mass to be determined based on the velocities of galaxies. Quite simply, the virial theorem states that the average kinetic energy is equal to minus one-half times the average gravitational potential energy. In 1933, Zwicky measured the velocity dispersion of galaxies in the Coma cluster and found that objects in the periphery of the Coma cluster orbited faster than that expected from the mass as inferred by its luminosity [77]. With this discovery, Zwicky is usually credited as the first person to provide evidence for dark matter. Girardi et al has also applied the virial theorem on a sample of about 100 galaxy clusters from various surveys [78]. The masses calculated with the virial theorem agree well with the values determined by x-ray measurements. The determination of mass by x-rays from intergalactic gas is similar to methods used for elliptical galaxies. Many such studies have been performed [79, 80].

Dark matter has also been observed in galaxy clusters via gravitational lensing, which has the advantage of not requiring assumptions on the internal dynamics, unlike the previously described methods. There are several classes of lensing. In strong gravitational lensing, the effect is very visible resulting in arcs and multiple images of the background object. Such cases are few and difficult to analyze due to the high non-linearity of the effect. In weak lensing, the distortions (ellipticity vectors and brightness) are much more minute and can only be detected by analyzing a large number of sources. Gravitational lensing measurements have confirmed the results of virial and x-ray analyses [81]. An overview of weak lensing can be found in [82, 83].

Data from the Sloan Digital Sky Survey (SDSS) has been used to probe the halo mass profile of isolated galaxies. This study showed halo profiles agreeing with that of dark matter models and contradicting those of MOND [84]. A flat rotation curve (which corresponds to a density distribution of $\rho \propto r^{-2}$) agrees with MOND while a falling velocity dispersion ($\rho \propto r^{-3}$) is predicted by galactic simulations of dark matter halos. The key to distinguishing
between halo profiles is to measure the velocities of objects at very large radii (<100 kpc). Tracers such as neutral hydrogen (HI) and x-ray emissions from diffuse hot gases in isolated elliptical galaxies indicate the presence of dark matter halos but neither indicators extend far enough to distinguish halo profiles. Strong gravitational lensing also provides a probe of only the inner regions. Weak gravitational lensing has also failed to distinguish between models. Observation of the motion of satellites (smaller galaxies) of galaxies provide an excellent probe of the dark matter distribution at large radii but statistics had been limited until SDSS. In this study, about 3,000 satellites with absolute blue magnitudes extending down to $M_B = -14$ were found. The velocity dispersion of these satellites were found to decrease with distance from the primary. This was the first direct observational confirmation of density decline which agrees with all cosmological models. Furthermore, the results contradict alternative theories such as Modified Newtonian dynamics (MOND).

1.2.3 Evidence at the Cosmological Level

Cold non-baryonic dark matter is a key component to our best models of the universe. In these models, dark matter is needed to explain the anisotropies and polarization in the cosmic microwave background (CMB) and the formation of large scale structures. Also, some dark matter is needed to explain the abundance of primordial elements generated by Big Bang nucleosynthesis (BBN). Experimental confirmation of these models support the idea that most of dark matter in the universe is cold and non-baryonic.

Cosmic Microwave Background

The CMB was first proposed by Robert Dicke in 1946 and was first detected by Penzias and Wilson in 1964 [85]. Its detection gave much support for the Big Bang Theory. About 379,000 years after the beginning of the universe when the temperature had dropped to about 3000 K, the electrons combined with the protons to form hydrogen. At this point, the photons decoupled from matter as they were no longer energetic enough to ionize hydrogen. The temperature of the photons continued to decrease as the universe expanded. These are
Fig. 2.— The WMAP 5-year TT power spectrum along with recent results from the ACBAR experiment. The red curve is the best-fit ΛCDM model to the WMAP data, which agrees well with all data sets when extrapolated to higher-$\ell$. The other experiments calibrate with WMAP or WMAP’s measurement of Jupiter (CBI). The CMB photons. The CMB spectrum almost perfectly follows the theoretical curve of a $T = 2.725$ K black body and is isotropic up to roughly 1 part in 100,000. The CMB essentially gives a snapshot of the structure of the early universe, and dark matter, if present, would be apparent in the CMB as it affects the evolution of this structure.

There are several sources of anisotropies [86]. The CMB is maximally blue shifted in one direction and maximally redshifted in the opposite direction. This dipole anisotropy, which is roughly 1 part in 1000, is due to the movement of our frame of reference relative to the CMB rest frame. The other sources of anisotropies can be divided into two categories. The primary anisotropies are due to processes occurring until the time of recombination and the secondary fluctuations are due to phenomena occurring since then. Some basic mechanisms for the primary anisotropies are acoustic oscillations and diffusion damping (also known as collisional dissipation or Silk damping). The acoustic oscillations are due to
two opposing effects in the photon-baryon plasma. The baryons accumulate by gravitational force while the pressure in the photons pushes the gas apart. This results in oscillations which give the characteristic peak structure of the CMB. Diffusion damping occurred during recombination as the mean free path of photons increased. The photons diffused from hot regions to cooler ones while dragging along protons and electrons. This reduces smaller scale anisotropies. Some secondary anisotropy contributions include the Sachs-Wolfe effect which causes the wavelength of photons to be shifted by changing gravitational potentials and the Sunyaev-Zel’dovich effect, which is essentially a reverse Compton interaction where high energy electrons transfer energy to CMB photons.

The anisotropies were first measured by the Far Infrared Absolute Spectrophotometer (FIRAS) on the COBE satellite [87]. The Wilkinson Microwave Anisotropy Probe (WMAP), launched into space in 2001, provides the most accurate measurement of the CMB [2]. Other prominent experiments include the balloon-borne BOOMERanG [88] and MAXIMA [89] and ground-based ACBAR [90] and DASI [91]. The Planck telescope, launching in April 2009, will improve upon the results of WMAP [92].

The angular power spectrum is obtained by decomposing the map into spherical harmonics, while taking into consideration various distortions such as emissions from galaxies. Figure 1.2 shows the power spectrum of the cosmic microwave background radiation temperature as measured by WMAP and other detectors. The size and location of the peaks provide valuable information on cosmological parameters, such as the curvature and matter/energy composition of the universe. WMAP provides a very precise composition of the universe: \( \Omega_\Lambda = 0.742 \pm 0.030 \) (dark energy), \( \Omega_c h^2 = 0.1099 \pm 0.0062 \) (cold dark matter) and \( \Omega_b h^2 = 0.02273 \pm 0.00624\% \) (baryonic) [93]. For \( h = 0.719 \), \( \Omega_c = 0.213 \) and \( \Omega_b = 0.0440 \).

**Big Bang Nucleosynthesis**

The theory of Big Bang nucleosynthesis, which was first proposed by Alpher and Gamov [94], describes the formation of light elements during the very hot phase of the universe. Matching the observed abundances of these elements to the predicted values constrains
20. Big-Bang nucleosynthesis

Figure 20.1: The abundances of $^4$He, D, $^3$He and $^7$Li as predicted by the standard model of big-bang nucleosynthesis. Boxes indicate the observed light element abundances (smaller boxes: $2\sigma$ statistical errors; larger boxes: $\pm 2\sigma$ statistical and systematic errors). The narrow vertical band indicates the CMB measure of the cosmic baryon density. See full-color version on color pages at end of book.

20.2. Light Element Abundances

BBN theory predicts the universal abundances of D, $^3$He, $^4$He, and $^7$Li, which are essentially determined by $t \sim 180$ s. Abundances are however observed at much later
the relative abundance of baryons in the universe ($\Omega_b$), which can be compared to the independently determined values from the CMB anisotropy measurements.

BBN began three minutes after the Big Bang when the universe was cool enough for protons and neutrons to form and combine into atoms, and ended about seventeen minutes later. BBN predicts the primordial abundances to be roughly 75% H (hydrogen comprising 75% of the baryon mass, not 75% of the atoms), 25% $^4$He, 0.01% D ($^2$H) and $10^{-10}$ $^7$Li [3]. No elements beyond beryllium are created because of the absence of stable nuclei with 5 or 8 nucleons (no nucleosynthesis via p$^4$He, n$^4$He or $^4$He$^4$He reactions). Figure 1.3 shows the measured and predicted abundances of $^4$He, D, $^3$He and $^7$Li relative to H as a function of the photon/baryon fraction ($\eta$) and the reduced baryon density ($\Omega_B h^2$). The evolution of the chemical abundances follow simple thermodynamics with consideration for the expansion of the universe. As the temperature decreased, the rate of conversion between protons and neutrons dropped faster than the Hubble expansion rate which resulted in a “freeze-out” of a non-equilibrium population. The freeze-out locked in a baryon population of 7 protons for every 1 neutron, which explains the observation of $^4$He comprising 25% of the mass – out of every 16 nucleons (2 neutrons and 14 protons), four of them combine to form $^4$He and thus make up 1/4 of the mass. This percentage is not strongly dependent on the initial conditions. BBN (or some other explanation) is needed to explain the amount of $^4$He because the observed amount is far larger than what could have been generated by stellar nucleosynthesis. The primordial abundance of $^4$He is inferred from H II regions within dwarf galaxies [95].

Unlike $^4$He, deuterium is not very stable and is easy to destroy. Two deuterium atoms can fuse into a $^4$He which is more stable. Not all of the deuterium is destroyed during BBN because of cooling. Its primordial abundance depends strongly on the baryon density as this dictates how much is destroyed. Since no post Big Bang processes are known to create deuterium and it is easily destroyed, most measured values can be seen as lower limits. Measurements of distant quasars provide the best estimate of its primordial abundance [96, 97]. The abundance of $^3$He is difficult to measure in extra-galactic sources because
its emission spectrum is more or less the same as $^4\text{He}$. Thus its proportion is deduced only from sources within our solar system and H II regions in our galaxy. Due to the difficulty in measuring its abundance and the uncertainty in stellar contribution [98], $^3\text{He}$ is not used as a probe. The amount of $^7\text{Li}$ can be deduced by probing the outermost layers of old stars where much less stellar nucleosynthesis has occurred and thus have a mixture that resembles the primordial fluid [99]. The predicted and measured $^7\text{Li}$ abundances are the same order of magnitude. The lack of a better agreement is more likely due to our incomplete understanding of stellar physics rather than of BBN.

Figure 1.3 shows the constraints by observed abundances that set a bound in the baryon fractional density. The concordance gives the baryon content as $0.017 \geq \Omega_B h^2 \geq 0.024$ (95% CL) [3] (or $0.033 \geq \Omega_B \geq 0.046$ for $h = 0.719$). Also shown in Figure 1.3 is the bound independently determined by WMAP measurements which remarkably agrees with the bounds set by BBN.

**Large Scale Structure**

The observed large-scale structure of the universe supports the existence of cold dark matter. The matter density distribution which is imprinted in the CMB should also be apparent in the current distribution of galaxies. The distribution has been determined by extensive redshift surveys such as the 2dF Galaxy Redshift Survey (2dFGRS) and Sloan Digital Sky Survey (SDSS). Figure 1.4 shows the reconstructed position of galaxies from the SDSS data. Over a million galaxies have been cataloged by these studies. The organization of structures beyond the size of super cluster was not well known until the discovery of structures such as the “Great Wall” by Geller and Huchra [100]. It is believed that the structures evolved from the “bottom up” – that is, smaller structures such as galaxies formed first, then clusters and super clusters. This hypothesis is supported by the age of stars and by the existence of currently forming super clusters. As simulations have shown, a universe with only baryons cannot lead to the structure we see today [101]. The baryons were too strongly coupled to the photons until recombination which is too late to grow these
structures. Non-baryonic cold dark matter allows structures to begin forming sooner as they are decoupled from the photons. A universe with hot dark matter would not have started galaxy formation soon enough and would have created structures that are too large [102]. Detailed analysis of the redshift survey data has yielded matter densities in accordance to that determined by WMAP. Analysis of the 2dFGRS data gives $\Omega_m = 0.26 \pm 0.05$ and $\Omega_b = 0.044 \pm 0.016$ assuming $h = 0.07$ based on a sample of 221,000 galaxies [103]. SDSS gives values in agreement [104]. Combining redshift survey data with WMAP results in significant improvements in cosmological parameters [105, 106, 107].

As shown in this section, we have discovered much evidence dark matter at various scales. The more sophisticated measurements at the cosmological scale strongly and independently
suggest that most of the dark matter is cold, non-baryonic and comprises about a quarter of the mass/energy of the universe. Most impressive is that observations of BBN, CMB and large scale structure formation are probes of the universe at vastly different times – the universe at several minutes, 379,000 years, and billions of years, respectively. In the next section, we review the dark matter candidates and alternative explanations.

1.3 Solutions to the Dark Matter Mystery

1.3.1 Baryonic Dark Matter

As discussed in the previous section, most of the matter in the universe is dark and non-baryonic. However, not all of the baryonic matter has been accounted for. The amount of luminous mass in the universe is $\Omega_{\text{lum}} \sim 0.005$ while $\Omega_b = 0.04$. A proposed candidate for baryonic dark matter is neutral hydrogen which is seen in intergalactic clouds at high redshifts ($z > 1$). If the density of this gas is low enough, these clouds might not have collapsed yet to produce stars. Neutral hydrogen is detected by observing the Lyman-\(\alpha\) forest (the absorption spectra of high-z quasar emissions). Measurements give only a lower bound of $\Omega_b h^2 \geq 0.018$ due to various systematics [108]. The Lyman-\(\alpha\) forest disappears at $z < 1$ – the dark baryons must have taken some other form. The most popular form has been the massive compact halo objects (MACHO), which includes dim objects such as brown dwarfs, black holes, white dwarfs, and neutron stars. These objects are detected indirectly through gravitational microlensing. Such searches involve monitoring numerous stars for a brief increase in light intensity when a MACHO crosses the direct line of sight between the observer and the star. Several microlensing searches for these objects have so far found too few incidents to account for all of the missing mass [109, 110]. In fact, MACHOs comprise at most 25% of the baryonic dark matter.
1.3.2 Non-Baryonic Dark Matter

Neutrino

Although neutrinos are stable and massive [111, 112], they are not good candidates for dark matter. As previously stated, even small amounts of neutrinos over dampen the small scale fluctuations in the CMB power spectrum. This happens because the neutrinos travel at such a high speeds that they smooth out any fluctuations in the matter density until they are cooled down by expansion. Also, it is difficult to see how neutrinos could populate the dark galactic halos [113]. WMAP and SDSS data have constrained neutrinos to comprise no more than 12% of dark matter [105].

Axions

The missing mass problem might be evidence of particles from extensions to the Standard Model. One such particle is the axion which is a low-mass neutral pseudoscalar particle that was first postulated by the Peccei-Quinn theory as a solution to the problem of strong CP violation in QCD [114, 115, 116]. Experiments and astrophysical arguments limit the mass to between $1 \times 10^{-6}$ and $1 \times 10^{-3}$ eV [117, 118, 119]. A summary of past and present searches can be found in [120]. Since axions are chargeless and have small cross sections for strong and weak interactions, they rarely interact with ordinary matter. Given their low mass, decay modes are unlikely which allows for the existence of a large non-thermal and cold relic population that “froze out” shortly after the Big Bang. A discussion on relic populations is given in the next section in the context of weakly interacting massive particles.

Axions couple to photons according to

$$\mathcal{L} = -\frac{1}{4} g_{a\gamma} F_{\mu\nu} F^{\mu\nu} a = g_{a\gamma} E \cdot B \ a$$

(1.2)

where $F$ is the electromagnetic field-strength tensor, $E$ and $B$ are the electric and magnetic fields, $g$ is the axion-photon coupling strength, and $a$ is the axion field. Several current experiments exploit this coupling and have a realistic chance for detection if they exist.
One such experiment is the Axion Dark Matter Experiment (ADMX) which exploits the conversion of axions into photons in sensitive microwave cavities [121, 122]. The expected signature of the axion signal is a peak of excess power at a frequency of $f = mc^2/h$, where $m$ is the mass of the axion, $c$ is the speed of light, and $h$ is the Planck constant. Since the axion mass is unknown, the microwave resonator cavity must be tunable. This experiment has excluded a local dark matter halo of KSVZ axions of mass $1.9 - 3.3 \times 10^{-6}$ eV with greater than 90% confidence [123].

If axions are real, we expect the sun to be a detectable source of axions. Solar axions are produced when x-rays scatter off electrons and protons in the presence of a strong electric field and are expected to have a broad energy spectrum with an average energy of 4.2 keV [124]. The CERN Axion Solar Telescope (CAST) searches for these axions with a detector that exploits the reverse Primakoff effect whereby a strong transverse 9.0 T magnetic field converts the solar axions back into x-rays of the same energy as the parent axion. Their first run has set a 95% CL for the axion mass at $m \leq 2 \times 10^{-2}$ eV [125].

The PVLAS experiment fires polarized light through a long vacuum region with a 5.5 T magnetic field and searches for anomalous rotations of polarization [126]. According to the theory, the vacuum becomes birefringent – photons with polarization aligned with the magnetic field are delayed as they are preferentially transformed into axions which travel slower than the speed of light. The PVLAS collaboration initially claimed detection of an irregular rotation corresponding an axion mass of 1–1.5 meV, but retracted their results upon obtaining a null result after upgrades [127]. Note that detection of axions in CAST and PVLAS does not give the axion background density.

**Weakly Interacting Massive Particles**

Another cold, stable, non-baryonic dark matter candidate is the weakly interacting massive particle (WIMP) which participates in only weak and gravitational interactions and has a mass of at least several tens of GeV (otherwise they would have already been observed in colliders). No known particle has these characteristics. Supersymmetry predicts the
Using the above relations ($H = 1.66 g^{-1} T^2$ Imp) and the freezeout condition $T = n_v/c v = H$, we find $(n_{XLS})_0 = \frac{(n_{XLS})_f}{100 (m_X/GeV) (\langle \sigma A v \rangle/10^{-27} \text{cm}^3 \text{s}^{-1})}$, \(\text{(3.3)}\)

where the subscript $f$ denotes the value at freezeout and the subscript 0 denotes the value today.

The current entropy density is $S_0 \sim 4000 \text{ cm}^{-3}$, and the critical density today is $\rho_c \sim 10^{-5} h^2 \text{ GeV cm}^{-3}$, where $h$ is the Hubble constant in units of $100 \text{ km s}^{-1} \text{ Mpc}^{-1}$, so the present mass density in units of the critical density is given by

$$\rho_{X} = \frac{\rho_X}{\rho_c}$$

The result is independent of the mass of the WIMP (except for logarithmic corrections), and is inversely proportional to its annihilation cross section.

Fig. 4 shows numerical solutions to the Boltzmann equation. The equilibrium (solid line) and actual (dashed lines) abundances per co-moving volume are plotted as a function of $x \equiv m_X/T$

Fig. 4. Comoving number density of a WIMP in the early Universe. The dashed curves are the actual abundance, and the solid curve is the equilibrium abundance. From [31].

existence of WIMP-like particles [5]. In this framework, every standard particle has a corresponding “superpartner”, with all quantum numbers being identical except for spin which differs by 1/2. The supersymmetric partners of the Z-boson, photon and neutral Higgs have the same quantum numbers and thus can mix to form the “neutralinos.” In many models, the lightest of the four neutralinos is the lightest supersymmetric particle (LSP), and if R-parity is conserved, this particle is a possible stable relic of the Big Bang. Its mass is not well constrained. Theorists favor a mass ranging from several GeVs to a few hundred TeVs. The LEP accelerator has set a $\sim 40 \text{ GeV}$ lower limit [121]. To preserve unitarity, an upper limit of $3.2 \text{ TeV}$ has been suggested [128].
If the WIMP (which we refer to as $\chi$) exists and is stable, a cosmological relic of WIMPs could have been produced during the Big Bang [129, 130]. At $T > M_\chi$, the $\chi$’s would have existed in thermal equilibrium, decaying to lighter particles and vice versa. After $T < M_\chi$, the $\chi$’s annihilate ($\chi\bar{\chi} \rightarrow ll$, assuming that $\chi$ is a Majorana particle) and drop rapidly in abundance. When the annihilation rate drops below expansion, the population “freezes out,” leaving behind the present population. This idea has been used to constrain the mass of the heavy neutrino [131].

The relic abundance can be approximated by assuming an energy independent WIMP-annihilation cross section. This calculation and a more rigorous one are described in [5]. The number density of $\chi$ at thermal equilibrium is

$$n^\text{eq}_\chi = \frac{g}{(2\pi)^3} \int f(p) d^3p,$$

(1.3)

where $g$ is the number of internal degrees of freedom of the particle and $f(p)$ is the Fermi-Dirac or Bose-Einstein distribution. At $T \gg m_\chi$, $n^\text{eq}_\chi \propto T^3$ while at $T \ll m_\chi$, $n^\text{eq}_\chi \approx g(m_\chi T/2\pi)^{3/2} \exp(-m_\chi/T)$ – the WIMP abundance drops exponentially at a rate of $\Gamma\langle\sigma_A v\rangle n_\chi$ where $\langle\sigma_A v\rangle$ is the thermally averaged total annihilation cross section multiplied by the relative velocity. Freeze-out occurs when $\Gamma < H$, where $H$ is the Hubble constant. At this point the density of WIMPs is too low in density for annihilation to continue. The Boltzmann equation describing the evolution of the WIMP density is

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle\sigma_A v\rangle [(n_\chi)^2 - (n^\text{eq}_\chi)^2].$$

(1.4)

The second term on the left hand side represents the effects of expansion. The first and second terms account for the annihilation and creation of WIMPs, respectively. This equation applies for Dirac and Majorana particles. The Hubble expansion term is negligible at earlier times (and the density follows the thermal equilibrium abundance) and dominates at later times. The freeze out temperature is provided by $\Gamma(T_f) = H(T_f)$ which gives $T_f \approx m_\chi/20$. The Hubble expansion rate falls with temperature as $H(T) = 1.66g^*_s T^2/m_{Pl}$ where $m_{Pl} = 1.22 \times 10^{19}$ GeV and $g_s$ is roughly equal to the number of bosonic relativistic degrees of freedom plus 7/8 times the number of fermionic relativistic degrees of freedom.
The \( \chi \) mass density as a fraction of the critical density, \( \Omega_\chi \), can be calculated as a function of the annihilation rate. Note that the entropy per unit comoving volume \( s = 0.4 g_s T^3 \) is constant and thus \( n_\chi/s \) is constant as well. This gives

\[
\frac{(n_\chi/s)_0}{(n_\chi/s)_f} = \frac{H(T_f)}{\langle \sigma_A v \rangle s(T_f)}
\]

\[
= 1.66 \frac{1}{0.4 g_s^{1/2} T_f m_P \langle \sigma_A v \rangle}
\]

\[
\approx 100 \frac{1}{g_s^{1/2} m_\chi m_P \langle \sigma_A v \rangle},
\]

where the subscripts 0 and \( f \) denote values for today and at freeze-out, respectively. The entropy density today is \( s_0 \approx 4000 \text{ cm}^{-3} \) and the critical density is \( \rho_c \approx 10^{-5} h^2 \text{ GeV cm}^{-3} \).

The mass density of WIMPs is

\[
\Omega_\chi h^2 = \frac{m_\chi n_\chi}{\rho_c}
\]

\[
\approx 3 \times 10^{-27} \text{ cm}^3 \text{s}^{-1} \frac{1}{\langle \sigma_A v \rangle},
\]

which is independent of \( m_\chi \). As expected, the mass density is inversely proportional to the annihilation cross section. Figure 1.5 shows the equilibrium density and comoving abundances as a function of inverse temperature for various annihilation cross sections.

We have derived the relic density that a stable particle should have if it exists. The annihilation cross section of a new particle interacting at the weak scale can be estimated as \( \langle \sigma_A v \rangle \sim \alpha^2 (100 \text{ GeV})^{-2} \sim 10^{-25} \text{ cm}^3 \text{s}^{-1} \) where \( \alpha \sim 10^{-2} \). This value is close to that projected by cosmological arguments which strongly suggests that if a stable particle associated with the electro-weak scale interactions exists then it is likely to be the dark matter particle. This coincidence has provided strong motivation for finding WIMPs. The rich and diverse community of experimental searches is described in the next chapter.

### 1.3.3 Modified Gravity

Modified Newtonian dynamics (MOND) was first proposed by Mordehai Milgrom in 1981 as an explanation to the galaxy rotation problem [132]. Milgrom proposed a modification
of Newton’s second law of motion: at very small accelerations \((a < 1.2 \times 10^{-10} \text{ m s}^{-2})\), the gravitation force is proportional to \(a^2\) and not \(a\) and thus the velocity falls off with \(r^{-1}\) instead of \(r^{-2}\). This gives a flat velocity curve far from the center. However, it has been shown that 80\% of the mass in galaxy clusters is unaccounted for in a MOND framework [133] and as previously stated, observations of satellites orbiting about isolated galaxies from SDSS give a halo profile that contradicts that predicted by MOND [84]. Also, MOND is not a relativistic formulation of gravity so gravitational waves, gravitational lensing, cosmological expansion and other standard phenomena derived from general relativity are notably absent. In response to these deficiencies, Bekenstein proposed a relativistic theory called Tensor-Vector-Scalar (TeVeS) that reduces to MOND in the Newtonian regime [134]. This theory also appears to produce gravitational lensing. Improved measurements of the third acoustic peak of the CMB anisotropies places severe restrictions on these theories [135] in strong favor of the cold dark matter model. Also, these theories do not offer an adequate explanation for the gravitational lensing observations of colliding galaxy clusters where the dark matter separates from the normal matter. The hot gases of the two clusters interact electromagnetically and slow down while the stars and dark matter passed through. Two such examples are the Bullet cluster (1E 0657-56) [136] and the MACS J0025.4-1222 [137].
Chapter 2

Dark Matter Searches

2.1 Introduction

The great importance of detecting WIMPs has given birth to many experimental searches. Thus far, no uncontested experimental evidence for WIMPs has emerged. These experiments can be divided into three categories: indirect detection, direct detection and high energy collider searches. This chapter gives an overview of the searches.

2.2 Indirect Detection Experiments

Indirect detection experiments search for the decay products of WIMP-pair annihilations. Since a pair annihilation requires two WIMP particles, the rate scales with $n^2$ and thus it is best to examine potential wells such as the Earth and Sun, where they can accumulate through inelastic collisions. Since WIMPs are slow moving, they decay into products with an energy of $M_\chi$ and if they are massive enough, they may annihilate into gauge bosons and charged Higgs. We can not directly observe these particles but we can potentially detect their decay products. For the decay products to be observable, they must be stable and be distinguishable from background. Four candidates meet these requirements: gamma rays, neutrinos, anti-protons and positrons.
2.2.1 Gamma Ray Experiments

WIMP annihilations generate gamma rays via $\chi\chi \rightarrow \gamma\gamma$ and $\chi\chi \rightarrow \gamma Z$ with energies of $M_\chi$ and $M_\chi(1 - M_Z^2/4M_\chi^2)$, respectively and since the WIMPs are slow moving, these gammas are mono-energetic. As WIMPs, by definition, do not couple directly to photons, the Feynman diagrams for these two annihilation channels include loops (see Figure 2.1) and thus are suppressed. In general, the photon flux from WIMP annihilation is

$$\phi(E, \Delta\Omega) \sim \frac{\sigma_v}{M_\chi^2} \times \int_{\text{l.o.s}} \int_{\Delta\Omega} \rho^2(s) \, ds \, d\Omega$$

where $\sigma$ is the WIMP annihilation cross section, $v$ is the relative speed of two interacting WIMPs in the center of mass frame, $\rho$ is the mass density, and $M_\chi$ is the mass of the WIMP. The integral expression integrates the WIMP density over the line of sight and solid angle. The density of the dark halo generally takes the form of

$$\rho(r) = \frac{\rho_c}{(r/a)^\alpha \times (1 + (r/a)^\gamma)^{\beta - \gamma}/\alpha},$$

where $r$ is the distance from the center of the galaxy, $a$ is a model dependent radius of the galactic core, and $\rho_c$ is a constant that normalizes the dark matter density to that of the local density (0.3–0.5 GeV cm$^{-3}$) [138]. The parameters, $\alpha$, $\beta$, and $\gamma$ depend on the model and have ranges of 1–2, 2–3 and 0–1.5, respectively [138]. The optimal source is nearby, has a large concentration of WIMPs, and has a low background. The galactic center has an enhancement due to the higher density of WIMPs but also has a larger background. The galactic halo, on the other hand, has a lower background but also has a lower signal due to the lower WIMP density.

Space telescopes allow for the direct detection of these high energy gamma rays. These detectors typically consist of an anti-coincidence shield, a layer of high-Z material, a spark chamber or silicon particle tracker, and a NaI or CsI scintillation calorimeter (see Figure 2.2). Gammas are detected by watching for pair conversions to electrons and positions. Reconstructing the energies and tracks of the charged particles gives the gamma energy and direction. The anti-coincidence shield for removing charged particle background is...
When evaluated numerically, it is found that annihilation to the two-gluon final state can dominate annihilation to tree-level final states when the neutralino is very nearly pure gaugino (but lighter than the top quark), and it can be comparable when the neutralino is a mixed state. It is generally not significant when the neutralino is purely higgsino.

In addition to the two-gluon final state, the three-body \(qqg\) final state (which arises at lower order in perturbation theory) should also be considered \[252, 261\]. It seems to be the case, however, that the \(qqg\) final state may be more important than the \(gg\) final state only in regions of parameter space where the \(gg\) final state is itself unimportant. Therefore, the \(qqg\) final state can safely be ignored in most models \[252\].

6.6. Photon final states

Annihilation to photon pairs may have interesting observational consequences \[263-266\]. and will be further discussed in Section 10.3. This annihilation process is the most complicated of the two-body annihilation channels, possessing all the structure of the two-gluon channel plus several extra contributions.

The diagrams for the process \(\chi \chi \rightarrow \gamma \gamma\) are shown in Fig. 18. Annihilation proceeds through sfermion–fermion and charged–Higgs–chargino loops, which are similar to those that appear in the two-gluon amplitude (see Fig. 17), and by chargino–W-boson loops. The calculation of the amplitude for annihilation to photon pairs has a long history, and there exist several partial calculations in the literature. The calculation of Ref. [263] was carried out assuming that the neutralino was a pure photino or higgsino state, and in the limit of large sfermion masses. The calculation of Ref. [265] was also carried out in the limit of large sfermion masses but with

![Diagrams](https://via.placeholder.com/150)

Fig. 18. Diagrams contributing to neutralino annihilation into photons: (a) fermion-sfermion loops, (b) charged-Higgs-chargino loops, and (c) chargino-W-boson loops. From Ref. [266].
Figure 2.2: Diagram of the gamma-ray telescope measurement technique used in EGRET and LAT. The gammas convert into an electron and positron in the tungsten strips. Their positions are sensed by the strip detector and their energies are measured by an NaI or CsI calorimeter. Surrounding the detector is a high efficiency anti-coincidence shield. This diagram is from [6].
usually a scintillator coated dome surrounding the rest of the detector. This background has an intensity that is $\sim 10^5$ greater than that of the gamma-rays. An overview of the potential of observing these gamma rays for dark matter can be found in [139].

The Energetic Gamma-Ray Experiment Telescope (EGRET) was one of four gamma ray detectors on the Compton Gamma Ray Observatory (CGRO), which ran from 1991 to 2000. This detector had a 0.5 sr field of view, an energy range from 20 MeV to 30 GeV and an energy resolution of 10%. EGRET measured the diffuse emission from the Galactic plane discovering that at energies above $\sim 1$ GeV, the flux was $\sim 60\%$ higher than the predicted value [140], an excess which suggests the presence of WIMP annihilation [141, 142]. From the spectral shape of the excess gamma-rays, the WIMP mass is estimated to be between 50 and 100 GeV. Others have claimed that the dark matter interpretation is ruled out by anti-proton fluxes [143] and that this observation is due an inaccurate estimation of the sensitivity at around 1 GeV [144]. This issue may be resolved with the Large Area Telescope (LAT) aboard the Gamma-ray Large Area Space Telescope (GLAST) [145], which was launched into orbit in June 2008. GLAST will observe a larger energy range (20 MeV to 300 GeV), greater field of view (20% of the sky), and larger collection area (8000 cm$^2$ versus EGRET’s 1500 cm$^2$) with a resolution of several arcminutes for the highest energy photons and about 3 degrees for 100 MeV gammas. The large energy range will provide some overlap with ground based detectors.

While restrictions on the size and flux will always limit the maximum energy that space telescopes can observe, ground-based experiments do not have such limitations. These experiments have been observing the night sky for Cerenkov light emitted by gamma rays interacting with the upper atmosphere. These bursts of light have a duration of 3–4 ns and a wavelength of 300–400 nm, and typically cover an area of land with a radius on the order of 100 m. These detectors, called Imaging Atmospheric Cerenkov Telescopes (IACTs), have an energy range of 100 GeV to several TeVs and usually examine a region at $\sim 10$ km high where the showers reaches their maximum intensity. The telescopes typically consist of a large number of mirrors pointing toward a camera consisting of many PMTs.
Figure 2.3: Comparison between gamma (left) and cosmic ray (right) induced showers, the latter of which tends to be broader, less smooth and occur closer to the surface of the earth.
The three secondary mirrors at the 49 m level are spherical with a nominal diameter of 1.9 m and a focal length of 2.0 m. Each is composed of seven identical hexagonal facets made from front-surfaced aluminized glass in order to retain a high reflectivity at ultraviolet wavelengths, where most of the Cherenkov light from air showers is produced. The two secondary mirrors at the 37 m level are single spherical mirrors with a diameter of 1.1 m and a focal length of 1.1 m.

The secondary mirrors focus the light from the heliostats, which arrives as a wide beam, onto photomultiplier tube (PMT) canisters anywhere laterally on the slotted plate and to a slotted plate. With this system, it is possible to position the PMTs anywhere laterally on the slotted plate and to a slotted plate. The secondary mirrors focus the light from the heliostats, which arrives as a wide beam, onto photomultiplier tube (PMT) canisters anywhere laterally on the slotted plate and to a slotted plate. With this system, it is possible to position the PMTs anywhere laterally on the slotted plate and to a slotted plate.

C. Cameras

The final stage in the STACEE optics chain is the camera. There is one camera for each secondary mirror. The cameras at the 37 m and 49 m levels have a field of view of 40° and 20°, respectively. A PMT camera consists of a PMT and light concentrator enclosed in a high-reflectivity, high-spatial-frequency (HSF) reflector. The PMT has a diameter of 1.1 m and a focal length of 1.1 m.

The cameras at the 49 m level consist of 16 PMT assemblies and the cameras at the 37 m level consist of 8 PMT assemblies. Each PMT assembly consists of a PMT and light concentrator enclosed in a high-reflectivity, high-spatial-frequency (HSF) reflector. The PMT has a diameter of 1.1 m and a focal length of 1.1 m.

The secondary mirrors at the 37 m level are single spherical mirrors with a diameter of 1.1 m and a focal length of 1.1 m. Each is composed of seven identical hexagonal facets made from front-surfaced aluminized glass in order to retain a high reflectivity at ultraviolet wavelengths, where most of the Cherenkov light from air showers is produced. The two secondary mirrors at the 37 m level are single spherical mirrors with a diameter of 1.1 m and a focal length of 1.1 m.

The secondary mirrors focus the light from the heliostats, which arrives as a wide beam, onto photomultiplier tube (PMT) canisters anywhere laterally on the slotted plate and to a slotted plate. With this system, it is possible to position the PMTs anywhere laterally on the slotted plate and to a slotted plate. The secondary mirrors focus the light from the heliostats, which arrives as a wide beam, onto photomultiplier tube (PMT) canisters anywhere laterally on the slotted plate and to a slotted plate. With this system, it is possible to position the PMTs anywhere laterally on the slotted plate and to a slotted plate.

D. Photomultiplier Tubes

The PMTs are supplied with high voltage from LeCroy high voltage power supplies, which are controlled by a microcontroller. The PMTs are supplied with high voltage from LeCroy high voltage power supplies, which are controlled by a microcontroller. The PMTs are supplied with high voltage from LeCroy high voltage power supplies, which are controlled by a microcontroller. The PMTs are supplied with high voltage from LeCroy high voltage power supplies, which are controlled by a microcontroller. The PMTs are supplied with high voltage from LeCroy high voltage power supplies, which are controlled by a microcontroller.

E. Front-End Electronics

The front-end electronics include high voltage supplies, pulse processing, and trigger logic. The high voltage supplies are controlled by a microcontroller. The pulse processing includes a charge-to-pulse converter, a pulse整形器, and a pulse height analyzer. The trigger logic includes a trigger decision system and a trigger discriminator. The front-end electronics include high voltage supplies, pulse processing, and trigger logic. The high voltage supplies are controlled by a microcontroller. The pulse processing includes a charge-to-pulse converter, a pulse整形器, and a pulse height analyzer. The trigger logic includes a trigger decision system and a trigger discriminator.
two orders of magnitude too low for observing the predicted flux. Gamma detection of spiral galaxies such as M31 are difficult to interpret in terms of dark matter annihilations because of the presence of non-thermal processes, cosmic ray acceleration, and other poorly understood physical mechanisms. Bad weather prevented CELESTE from observing the Draco spheroidal galaxy, which is a better source as it is nearby and dominated by dark matter. This source was observed by the STACEE experiment which used the mirrors of the National Solar Thermal Test Facility (NSTTF) in Albuquerque, New Mexico [155, 156]. No excess in the gamma-ray signal was observed from Draco.

The Milagro experiment uses a large 5000 m$^2$ pool of water as a target for gamma ray induced showers. PMTs in the water detect the Cerenkov light emitted by the secondaries of the air shower. Milagro has a wide view but a low efficiency for background rejection. This detector has observed the sun for neutralino annihilations, finding no statistically significant signal [157].

2.2.2 Neutrino Experiments

WIMP annihilations can release neutrinos. Since neutrinos interact so rarely with normal matter, very massive natural structures such as the ocean and polar ice caps are needed. Muon neutrinos interact with target via W-boson exchange generating muons, which generates Cerenkov light. Electron neutrinos are not seen because resulting electrons are quickly absorbed before reaching detectors. Tau neutrinos are not detected because the tau generation is suppressed by its high mass. The neutrino interaction cross section is proportional to $E_\nu$ and the range of the muon is proportional to $E_\mu$, so there is an $E_\nu^2$ enhancement, thus favoring detectors with high energy ranges. A continuous spectrum (rather than a line) is expected as neutrinos are produced not only through $\chi\chi \rightarrow \nu\bar{\nu}$ but also via $\chi\chi \rightarrow f\bar{f}$, where $f$ may decay into a neutrino or hadronize and decay into a neutrino if it is a quark. Neutrinos from the sun can also have their energy shifted as they escape. These detectors must also be able to reject atmospheric neutrinos which have a flux that is $10^6$ greater than that of the upward muons.
Some experiments which have searched for these neutrinos include Baksan [158], Super-Kamiokande [67], MACRO [159], and AMANDA [160]. None of these experiments have observed excess neutrinos but they have set flux upper limits. Larger experiments are being constructed. The ANTARES detector consists of 900 downward facing PMTs attached onto 12 strings at a depth of about 2500 m in the waters of the French Mediterranean coast near Toulon [161]. Light scatters less in water than in ice but has more background light from sources such as bioluminescent organisms and radioactive salts. The successor to the AMANDA detector is IceCube, which is similar to ANTARES but uses the ice of the South Pole [162]. When completed, IceCube will be the largest neutrino detector with 4800 PMTs covering a fiducial volume of 1 km$^3$. The PMTs are suspended in the fiducial volume by 1.5 mile long strings. IceCube views events arriving from below which filters out the otherwise overwhelming muon rate. Due to the low angle of the sun at the South Pole, the ability to separate solar neutrinos from background is weak. The IceCube experiment is likely to exclude portions of supersymmetric space that have yet to be omitted by direct searches.

There may also be an amplification of WIMP annihilations in the Galactic center which can be observed as an enhancement in the neutrino flux as shown by Gondolo and Silk [7]. Dark matter halos are expected to have a steep central cusp with a density profile following $\rho \propto r^{-\gamma}$. Semi-analytical calculations find $\gamma \sim 1–2$ and simulations give $\gamma = 0.3–1.5$. The presence of a large black hole of mass $\sim 2 \times 10^6 M_\odot$ at the galactic center would lead to a large “spike” in the annihilation products in the surrounding region. The evidence for the existence of such a super massive black hole is overwhelming with observations of nearby objects and gives a density profile, $\rho \propto r^{-\gamma}$, with $\gamma \geq 3/2$. Observation of the annihilation products in this spike can help us determine the type of halo core, as a central cusp would result in a magnification in annihilation. Observations favor a softened core density which lead to only a small enhancement in annihilation rates. If the halo has a central cusp, as expected from simulations, the enhancement can be five orders of magnitude or more. An upper limit to the muon flux can impose an upper bound on $\gamma$ or indicate that the neutralinos are not the dark matter particle annihilating there. Future experiments in the Northern
So we restrict following Bergström and Ahn's points in parameter space. We use 35,121 points for the cold dark matter annihilation density. The relic density comes from the annihilation requiring that the neutralino is a good cold dark matter candidate, in the sense that its relic density satisfies the bound. Figure 2.5 is a plot of the enhancement as a function of $\gamma$. Figure 2.6 shows the expected flux of neutrino-induced muons from the galactic center with and without the spike around the black hole. For kinematic reasons, the background flux of low energy antiprotons is suppressed. The idea of searching for excess antiprotons in cosmic rays was first proposed in the 1980s [164, 165].

A strong enhancement results instead for power-law profiles. We find

$$S_{\text{spike}} = \rho_D \sigma v D \left( \frac{R_{\text{sp}}}{D} \right)^{2-\gamma}$$

where $R_{\text{sp}}$ is given after Eq. (9) with $\rho_0$ replaced by $\rho_D$ and $r_0$ by $D$. We fix $R_{\text{in}}$ so as to match the integration of the numerically calculated density profile including capture and annihilation: we find that $R_{\text{in}} = 1.5 [(20R_p)^2 + R_{\text{sp}}^2]^{1/2}$ gives a good approximation to the flux (within 6% for our values of $\gamma$).

Contrary to the case of finite cores, for cusped halos, there is a huge increase in flux from the galactic center when the spike is included, typically 5 orders of magnitude or more, unless the inner halo slope $\gamma$ is very small.
be disregarded in dark matter studies. So we restrict the study of other signals. We focus on the neutrinos and postpone the study of other signals.

Contrary to the case of finite cores, for cusped halos there is a huge increase in flux from the galactic center when the spike is included, typically 5 orders of magnitude or more, unless the inner halo slope $\gamma$ is very small.
Given the uncertainties in production and propagation of secondary antiprotons, it is uncertain whether these measurements will further constrain WIMP models [8].

In the 1990s, it was proposed that if the WIMPs are heavier than $W^\pm$ and $Z^0$ bosons and have a significant annihilation branch into $W^\pm$ and $Z^0$ pairs, then the branches ($W^+ \rightarrow e^+ + \nu_e$ and $Z_0 \rightarrow e^+ + e^-$) should provide a very distinct feature in the cosmic-ray position spectrum (as seen in Figure 2.8) [9]. These positrons have an energy of roughly half of the neutralino mass. The predicted peaks may be washed out by boosting in the case of $m_\chi \gg m_{W,Z}$ or by energy loss as the positrons travel across the galaxy. The gauge bosons also decay into quarks and pions, which consequently decay to produce positrons. These positrons form another peak at $M_\chi/20$.

PAMELA has recently published preliminary results of its measurement of the cosmic ray positron fraction up to 100 GeV [10]. This measurement contained much higher statistics and probed higher energies than any previous measurement. Figure 2.9 compares PAMELA’s measurement with a calculation for secondary production of positrons during the propagation of cosmic ray nuclei. The observed rise may indicate annihilation of dark matter in the galactic halo or production of positrons by near-by pulsars [167].

Some indirect detection experiment may exclude portions of supersymmetric space not yet excluded by other experiments but unless astrophysical backgrounds are fully understood there will always be some uncertainty. In direct detection experiments, there is no ambiguity once backgrounds and other detector systematics are sufficiently suppressed or understood. However, even if dark matter is found first in direct detection experiments or colliders, these indirect detection experiments are still needed to probe the halo distribution.

2.3 Direct Detection Experiments

The race for direct detection of the WIMP began in the 1980s with the first experiments using conventional high purity germanium and silicon detectors at liquid nitrogen temperatures. In the next decade, NaI scintillation detectors gained favor as they have some background discrimination capabilities based on the scintillation pulse shape. DAMA is the
shown in Fig. 3, the antiproton spectrum, in which a characteristic peak at 2 GeV is dominated by the cosmic-ray component. The measured agreement of these tests of models of secondary production makes a priori improbable to attribute to the secondary component alone the measured proton and helium data on cosmic-ray spectra for the solar minimum period is expected to be very small [20].

Also shown in Fig. 3 is the secondary antiproton production parameter $\lambda(R,\beta)$ for $\beta = 10^\circ$ at solar minimum. The diffusion model [23,24], in which the ratio, under the hypothesis that the secondary antiproton for the solar minimum period is relevant to the cosmic-ray species (nuclear), can be represented as $\lambda(R,\beta)$, $\beta$ by propagation processes. Essential inputs required by which are similar to than previous data and are relevant to the $\lambda(R,\beta)$.

These calculations reproduce our spectrum at the peak region remarkably well, under the hypothesis that the different cosmic-ray production. The distinct peak structure at 2 GeV is dominated by the cosmic-ray component.
The continuum positron spectrum is then obtained by the convolution of the sources \( f_S(\epsilon) \) and \( f_\pi(\epsilon) \) with the Green's function [Eq. (2) or (4)]. The complete positron spectrum from decays of gauge bosons produced by halo annihilations of a Higgsino of mass 120 GeV is shown in Fig. 3. Note that the peak associated with the direct decays of \( W^+ \) and \( Z^0 \) bosons remains quite prominent and is not washed out by the continuum positron radiation. Moreover, there appears to be a second, less prominent peak centered at an energy of about \( m_\chi/20 \), which could

Since the composition of the Universe is of such great importance to both particle physics and cosmology, any and all avenues that can lead to the discovery of dark-matter candidates must be pursued. Here we have emphasized the distinctive feature in the cosmic-ray positron spectrum that arises in the Universe is of such great importance to both particle physics and cosmology, any and all avenues that can lead to the discovery of dark-matter candidates. Here we have emphasized the distinctive feature in the cosmic-ray positron spectrum that arises in the Universe is of such great importance to both particle physics and cosmology, any and all avenues that can lead to the discovery of dark-matter candidates.
Figure 2.9: Positron fraction as measured by PAMELA and theoretical prediction (black line) for pure secondary production of positrons during the production of cosmic-rays in the galaxy. Plot taken from [10].
Figure 2.10: Chart of direct detection searches categorized by measurement technique.

only group to have claimed detection of the WIMP. They have observed an annual modula-
tion using their 100 kg NaI detector [168]. However, many other experiments have since
contradicted these results. Also seen in the 1990s was the rise of sub-kelvin superconducting
semiconductor detectors which have excellent event-by-event discrimination through direct
measurement of phonons and ionization. By the mid-2000s, the same principle was demon-
strated in two-phase xenon and argon detectors, which have gained attention because of
their scalability. Also being developed are single-phase liquid neon and liquid argon detec-
tors with very powerful pulse shape discrimination. A comprehensive list of direct detection
experiments can be found in [169]. Figure 2.10 is a chart of the direct detection searches
categorized by measurement technique. In this section, we first give an overview of the
signal generation in such detectors and then provide a brief overview of past and current
experiments.

2.3.1 Physics of Direct Detection

Here, we give a brief summary of the calculation of the WIMP signal in direct detection
experiments. A more complete overview can be found in [170].
Interaction Rate

The WIMP-nucleon cross section upper limit is calculated by comparing the observed rate upper limit to that expected from the theoretical calculations. The expected differential event rate can be written as

\[ \frac{dR}{dE} = R_0 S(E) F^2(E) I, \]  

(2.3)

where \( R_0 \) is the total event rate; \( S(E) \) is the spectral function that includes the effects of orbital velocity of the earth, detection efficiency for nuclear recoils, and instrumental limitations; \( F \) is the form factor correction; and \( I \) is the factor related to the type of spin interaction. The rest of this section is describes these terms.

The differential dark matter particle density is

\[ dn = \frac{n_0}{k} f(v, v_E) d^3v, \]  

(2.4)

where \( k \) is a normalization constant, \( n_0 \) is the mean dark matter density, \( v \) is the velocity of the dark matter particle onto the target, and \( v_E \) is the velocity of the earth relative to the dark matter halo. The normalization constant, \( k \), is defined as

\[ k = \int_0^{2\pi} d\phi \int_{-1}^{1} d(\cos \theta) \int_{0}^{v_{esc}} f(v, v_E) v^2 dv, \]  

(2.5)

where \( v_{esc} \) is the escape velocity of the WIMP in the halo. The dark matter velocity distribution is usually taken to be Maxwellian:

\[ f(v, v_E) = \exp\left(-\frac{(v - v_E)^2}{v_0^2}\right). \]  

(2.6)

For \( v_{esc} = \infty \), \( k = k_0 = (\pi v_0^2)^{3/2} \) and with a truncation at \( |v + v_E| = v_{esc} \),

\[ k = k_1 = k_0 \left[ \text{erf}\left(\frac{v_{esc}}{v_0}\right) - \frac{2}{\sqrt{\pi}} \frac{v_{esc}}{v_0} e^{-v_{esc}^2/v_0^2}\right]. \]  

(2.7)

The event rate per unit mass of WIMPs traveling with velocity \( v \) onto a target of atomic mass \( A \) and nucleus cross section of \( \sigma \) is

\[ dR = \frac{N_0}{A} \sigma v dn, \]  

(2.8)
where \( N_0 \) is the Avogadro’s number \( (6.022 \times 10^{23}) \) and \( dn \) is the differential particle density defined in Equation 2.4. The total event rate is

\[
R = \frac{N_0}{A} \sigma_0 \int v \, dn \equiv \frac{N_0}{A} \sigma_0 n_0 \langle v \rangle,
\]  

(2.9)

where \( \sigma_0 \) is the zero-momentum cross section. The total event rate per unit mass for \( v_E = 0 \) and \( v_{esc} = \infty \) is defined as

\[
R_0 = \frac{2}{\sqrt{\pi}} \frac{N_0}{A} \rho_\chi \sigma_0 v_0,
\]  

(2.10)

where \( \rho_\chi \) is the halo dark matter density for which we will use a nominal value of 0.3 GeV \( c^{-2} \text{ cm}^{-3} \). By substituting \( R_0 \) into Equation 2.9, we obtain

\[
R = R_0 \frac{\sqrt{\pi}}{2} \frac{\langle v \rangle}{v_0}
\]  

(2.11)

\[
= R_0 \frac{k_0}{k} \frac{1}{2\pi v_0^2} \int v f(v, v_E) d^3 v,
\]  

(2.12)

and define the differential form as

\[
dR = R_0 \frac{k_0}{k} \frac{1}{2\pi v_0^2} v f(v, v_E) d^3 v.
\]  

(2.13)

\( R_0 \) can be written in units of \( \text{kg}^{-1} \text{ d}^{-1} \) or “tru.” With \( \rho_\chi = 0.3 \text{ GeV} \, c^{-2} \text{ cm}^{-3} \) and \( v_0 = 230 \text{ km s}^{-2} \), we obtain

\[
R_0 = \frac{377}{M_\chi M_T} \left( \frac{\sigma_0}{1 \text{ pb}} \right) \left( \frac{\rho_\chi}{0.3 \text{ GeV} \, c^{-2} \text{ cm}^{-3}} \right) \left( \frac{v_0}{230 \text{ km s}^{-1}} \right),
\]  

(2.14)

where \( M_T \) and \( M_\chi \) are in units of GeV \( c^{-2} \) and \( M_T = 0.932 \, A \).

The recoil energy of the target nucleus hit by a dark matter particle of energy \( E = \frac{1}{2} M_E v^2 \) and scattering at an angle of \( \theta \) is

\[
E_R = E \, r \, (1 - \cos \theta)/2,
\]  

(2.15)

where \( r = 4M_\chi M_T/(M_\chi + M_T)^2 \). Assuming isotropic scattering (the recoils are uniform in \( E_R \) between 0 and \( E_r \)), we obtain

\[
\frac{dR}{dE_R} = \int_{E_{\min}}^{E_{\max}} \frac{1}{E_r} dR(E)
\]  

(2.16)

\[
= \int_{v_{\min}}^{v_{\max}} \left( \frac{v_0}{v} \right)^2 dR(v),
\]  

(2.17)
where the $E_{\text{min}}$ ($E_{\text{max}}$) is the smallest (largest) energy that can give the recoil energy $E_R$; $E_0 = \frac{1}{2} M \chi v_0^2 = (v_0/v)^2 E$; and $v_{\text{min}}$ ($v_{\text{max}}$) is the velocity corresponding to the energy $E_{\text{min}}$ ($E_{\text{max}}$). Substituting in Equation 2.13, we obtain

$$
\frac{dR}{dE_R} = \frac{R_0}{E_0 r} \frac{k_0}{k_1} \frac{1}{2\pi v_0^2} \int_{v_{\text{min}}}^{v_{\text{max}}} \frac{1}{v} f(v, \chi v) d^3v. \quad (2.18)
$$

Integrating $v$ from 0 to $\infty$ for $v_E = 0$ gives the unmodified nuclear recoil spectrum

$$
\frac{dR(v_E = 0, \infty)}{dE_R} = \frac{R_0}{E_0 r} e^{-E_R/E_0r}, \quad (2.19)
$$

which shows that the dark matter recoil spectrum is a falling exponential. This explains the importance of having sensitivity at low energies. For other cases of $v_E$ and $v_{\text{esc}}$, the differential recoil spectrums are

$$
\frac{dR(0, v_{\text{esc}})}{dE_R} = \frac{k_0}{k_1} \frac{R_0}{E_0 r} \left( e^{-E_R/E_0r} - e^{-v_{\text{esc}}^2/v_0^2} \right), \quad (2.20)
$$

$$
\frac{dR(v_E, \infty)}{dE_R} = \frac{R_0 \sqrt{\pi} v_0}{E_0 r} \frac{v_E}{4} \left[ \text{erf} \left( \frac{v_{\text{min}} + v_E}{v_0} \right) - \text{erf} \left( \frac{v_{\text{min}} - v_E}{v_0} \right) \right], \quad (2.22)
$$

$$
\frac{dR(v_E, v_{\text{esc}})}{dE_R} = \frac{k_0}{k_1} \left[ \frac{dR(v_E, \infty)}{dE_R} - \frac{R_0}{E_0 r} e^{-v_{\text{esc}}^2/v_0^2} \right]. \quad (2.23)
$$

The spectrum $dR(v_E, \infty)/dE_R$ can be approximated as

$$
\frac{dR(v_E, \infty)}{dE_R} = c_1 \frac{R_0}{E_0 r} e^{-c_2 E_R/E_0r}, \quad (2.24)
$$

where $c_1$ and $c_2$ are fitting constants of order unity. The integral of Equation 2.24 gives the relation between the two constants. For most purposes, the average values of $c_1 = 0.751$ and $c_2 = 0.561$ can be used. The integrated rate of the recoil energy range between $E_1$ and $E_2$ is

$$
R(E_1, E_2) = R_0 \frac{c_1}{c_2} \left[ e^{-c_2 E_1/E_0r} - e^{-c_2 E_2/E_0r} \right]. \quad (2.25)
$$

**Form Factor Correction**

At higher momentum transfers ($q = \sqrt{2M \chi E_R}$), the de Broglie wavelength ($h/q$) becomes comparable to the size of the nucleus and the cross section begins to drop. The nuclear
Figure 2.11: Spin-independent nuclear form factor versus recoil energy (Equation 2.33, \( s = 1 \), and \( r_n = \sqrt{(5/3)(0.89A^{1/3} + 0.3)^2 - 5s^2} \) for various targets. The low form factor for xenon is the reason why xenon detectors must have a low energy threshold.

The form factor correction, \( F(q) \), is included to account for this effect. It is a function of the dimensionless quantity \( qr_n/h \), where \( r_n \) is the nuclear radius. With \( M_T = 0.932 \, A \), we have

\[
q(\text{MeV}^{-1}) = [2 \times 0.932(\text{GeV}^{-2})A \, E_R(\text{keV})]^{1/2},
\]

(2.26)

and with \( h = 197.3 \, \text{MeV} \, \text{fm} \) and \( r_n = a_n A^{1/3} + b_n \), we obtain

\[
qr_n(\text{dimensionless}) = 6.92 \times 10^{-3} A^{1/2} E_R^{1/2} (a_n A^{1/3} + b_n),
\]

(2.27)
where $E_R$ is in units of keV and $a_n$ and $b_n$ are in units of fm. The cross section can be separated into two parts as

$$\sigma(qr_n) = \sigma_0 F^2(qr_n),$$

(2.28)

where $\sigma_0$ is the zero-momentum cross section that contains information on the particular interaction, and $F$ is the form factor containing the momentum transfer dependence. The form factor is, in the first Born approximation, the Fourier transform of the density of scatterer centers, $\rho$:

$$F(q) = \int \rho(r) e^{iqr} d^3r$$

(2.29)

$$= \frac{4\pi}{q} \int_0^\infty r \sin(qr) \rho(r) dr.$$  

(2.30)

Assuming a thin shell density distribution gives

$$F(qr_n) = j_0(qr_n) = \frac{\sin(qr_n)}{qr_n},$$

(2.31)

where $j_0$ is the spherical Bessel function of index 0. For a solid sphere,

$$F(qr_n) = \frac{3j_1(qr_n)}{qr_n} = 3 \left( \frac{\sin(qr_n)}{(qr_n)^3} - \frac{\cos(qr_n)}{(qr_n)^2} \right).$$

(2.32)

For a more exact calculation of the form factor in the spin-independent case, we use a scatterer distribution that mirrors the charge distribution that is determined by electron and muon scattering. These distributions have a uniform distribution with a drop to zero over a “skin” thickness. One such distribution gives

$$F(qr_n) = 3j_1(qr_n) e^{-(qs)^2/2}$$

(2.33)

where $s$ is the skin thickness [171]. Figure 2.11 shows the spin-independent form factor (Equation 2.33) for various targets as a function of recoil energy.

In detailed calculations that include contributions from all nucleons, the form factor for the spin-dependent case has three parts representing the proton, neutron and interference terms or three terms representing the isoscalar ($p + n$), isovector ($p - n$) and interference terms. In the latter form, $F(qr_n)^2 = S(q)/S(0)$ where

$$S(q) = a_0^2 S_{00}(q) + a_1^2 S_{11}(q) + a_0 a_1 S_{01}(q),$$

(2.34)
Figure 2.12: Spin-dependent nuclear form factor (thin-shell approximation; Equation 2.31) versus recoil energy for various targets.

and $a_0$ and $a_1$ are the isoscalar and isovector coefficients, respectively, related to the WIMP-nucleon spin factors. The spin-dependent form factor (as seen in Figure 2.12) can be estimated with the thin shell approximation (Equation 2.31). We see that the form factors of spin-dependent and spin-independent interactions drops significantly with energy, which are additional reasons why having sensitivity to low energy recoils is so important.

Another component in the differential rate (Equation 2.3) is the spin component. The total zero-momentum WIMP-nucleus cross section can be written as

$$
\sigma_0 = 4G_F^2\mu_N^2 C,
$$

(2.35)

where $G_F$ is the Fermi coupling constant, $\mu_N = M_T M_\chi / (M_T + M_\chi)$, and $C$ is an en-
hancement factor that is different for spin-dependent ($C_{SD}$) and spin-independent ($C_{SI}$) interactions. For the spin-dependent interactions only unpaired nucleons contribute to the scattering amplitude, thus only nuclei with an odd number of protons and/or odd number of neutrons can have a spin-dependent cross-section. In the spin-dependent case,

$$C_{SD} = \frac{8}{\pi} \Lambda^2 J(J + 1),$$

(2.36)

where

$$\Lambda = \frac{1}{J} [a_p \langle S_p \rangle + a_n \langle S_n \rangle],$$

(2.37)

with $a_p$ and $a_n$ as the WIMP-proton and WIMP-neutron effective couplings, respectively, $\langle S_n \rangle$ and $\langle S_p \rangle$ as the neutron and proton spin expectation values of the nucleus, and $J$ as the total nuclear spin. More details can be found in [5]. In the spin-independent case,

$$C_{SI} = \frac{1}{\pi G_F^2} [Z f_p + (A - Z) f_n]^2,$$

(2.38)

where $f_p$ and $f_n$ are the effective WIMP-proton and WIMP-neutron couplings, respectively. For Majorana WIMPs, $f_p \approx f_n$ and

$$C_{SI} = \frac{1}{\pi G_F^2} (Af_p)^2.$$

(2.39)

This equation shows that the spin-independent cross sections follows $A^2$.

**Annual Modulation**

The motion of the Earth around the Sun induces an annual modulation in the WIMP signal. The velocity of the Earth is

$$v_E \approx v_s + v_c \sin \delta \cos(2\pi(t - t_0)),$$

(2.40)

where $v_s \approx 230$ km s$^{-1}$ is the velocity of the Sun (disc rotation), $v_c \approx 30$ km s$^{-1}$ is the rotational velocity of the Earth, $\delta = 30.7^\circ$ is the angle between the solar system axis and velocity vector of the Sun, and $t_0$ is June 2nd when the Earth’s speed is maximum [172]. Figure 2.13 is a plot of the average, June and December differential rates assuming
Figure 2.13: The average, June, and December differential rates. Also included is the difference between June (maximum) and December (minimum) rates. This assumes $\sigma_0 = 10^{-5}$ pb, $M_\chi = 100$ GeV $c^{-2}$, $v_0 = 230$ km s$^{-1}$, $v_s = 230$ km s$^{-1}$, $v_E = 30$ km s$^{-1}$, and $v_{esc} = 600$ km s$^{-1}$.

The absence of radioimpurities, low energy threshold ($\sim$1 keV), high energy resolution (1 keV) are some reasons why germanium is a good target material for dark matter detectors. The first Ge detectors were simply those built in the 1980s for detecting neutrinoless double...
beta decay of $^{76\text{Ge}}$. These first generation detectors helped eliminate Dirac neutrinos and cosmions [173] as possible dark matter candidates but were limited by electronics noise and the lack of background discrimination capability. Next generation ionization detectors use parts with less radioimpurities and separate modules for anti-coincidence background vetoing. Cosmogenically produced tritium is a concern because its 18.6 keV end-point results in background spanning the entire expected measurable WIMP spectrum. $^{68\text{Ge}}$ is another radioisotope that is cosmogenically generated from $^{70\text{Ge}}$. This problem can be partially alleviated by moving the detector underground or by using enriched high purity germanium (HPGe). The IGEX and H/M are two experiments which employ enriched $^{76\text{Ge}}$ for double-beta decay and WIMP detection [174]. IGEX-2002 has generated the best exclusion plot among detectors of this class [175]. The Heidelberg Dark Matter Search (HDMS) detector, running at Gran Sasso, consists of two $^{73\text{Ge}}$ enriched HPGe detectors, a p-type crystal surrounded by a well-type Ge crystal. The presence of $^{73\text{Ge}}$ is needed to probe the spin-dependent WIMP-nucleon cross section. This experiment has measured a dark matter limit with a final setup [176]. GEDEON plans to use an array of twenty-eight 2 kg crystals to reach a background rate of $10^{-3}$ dru below 100 keV with the expectation of exploring the annual modulation [177]. The GENIUS-TF is currently operating several detectors at Gran Sasso using natural abundance Ge detectors with a total mass of 10 kg [178]. The detector sits directly in liquid nitrogen, which serves as both a cooling medium and as a shield against external radioactivity. GERDA will first use the HPGe crystals previous used in H/M and IGEX and have a total mass of 15 kg and then move on to 20 kg of new modules. The detector will use water and liquid argon for shielding with the liquid argon also serving as a cooling medium [179]. The MAJORANA experiment will use two-hundred HPGe detectors with a total mass of 500 kg Ge, isotopically enriched to 86% in $^{73\text{Ge}}$ [180].
Figure 2: Model-independent residual rate of the single-hit scintillation events, measured by the new DAMA/LIBRA experiment in the (2 – 4), (2 – 5) and (2 – 6) keV energy intervals as a function of the time. The residuals measured by DAMA/NaI and already published in ref. [4, 5] are also shown. The zero of the time scale is January 1st of the first year of data taking of the former DAMA/NaI experiment. The experimental points present the errors as vertical bars and the associated time bin width as horizontal bars. The superimposed curves represent the cosinusoidal functions expected for a WIMP signal with the modulation amplitudes obtained by best fit over the whole data, that is: (0.0215 ± 0.0026) cpd/kg/keV, (0.0176 ± 0.0020) cpd/kg/keV and (0.0129 ± 0.0016) cpd/kg/keV for the (2 – 4) keV, for the (2 – 5) keV and for the (2 – 6) keV energy intervals, respectively. See text. The dashed vertical lines correspond to the maximum of the signal (June 2nd), while the dotted vertical lines correspond to the minimum. The total exposure is 0.82 ton×yr.

Figure 2.14: Model-independent residual rate of single-hit events in the energy ranges 2–4, 2–5, and 2–6 keV as a function of time over a 12 year period. Both the 0.29 ton-year DAMA/NaI and 0.53 ton-year DAMA/LIBRA data are included. The curves represent the cosinusoidal functions expected for the WIMP signal. The amplitudes obtained by best fit are (0.0215 ± 0.0026) cpd/kg/keV, (0.0176 ± 0.0020) cpd/kg/keV and (0.0129 ± 0.0016) cpd/kg/keV for the (2–4) keV, for the (2–5) keV and for the (2–6) keV energy intervals, respectively.
2.3.3 Solid Scintillation Detectors

In the mid-1990s, the performance of solid scintillation detectors eventually surpassed that of the Ge and Si detectors. Although these detectors have high gamma backgrounds, they are able to acquire greater sensitivity with pulse shape discrimination. Some of the early experiments include detectors built by BPRS [181], DAMA [181], ELEGANTS [182], Boulby DM collaboration [183]. These detectors generally have excellent long term stability allowing for the search of the annual modulation in the signal. The first direct detection experiment to report a positive annual modulation signal was DAMA/NaI in 1997 with a 3364 kg-day exposure in the winter and a 1185 kg-day exposure over the summer, using a detector consisting of nine 9.7 kg NaI detectors at Gran Sasso [184]. The results of additional exposures with this detector were also published again showing the annual modulation [168]. The DAMA collaboration recently published results from their DAMA/LIBRA (Large sodium Iodide Bulk for RARe processes) detector which consist of twenty-five 9.7 kg radiopure NaI modules [185, 11]. DAMA/LIBRA collected an exposure of 0.53 ton-year over four annual cycles which clearly displayed again the annual modulation in the single-hit residual rate. Figure 2.14 shows the residual rate in various energy bins of the DAMA/NaI and DAMA/LIBRA data as a function of time. Their observations, however, contradict the null results of many other experiments.

2.3.4 MilliKelvin Cryogenic Detectors

Motivated by the need for greater background discrimination power, the community explored cryogenic detectors that exploit the relative quenching in scintillation and charge signals of electron and nuclear recoils. These detectors usually run at millikelvin temperatures (using helium dilution refrigerators) and measure the energy through the collection of phonons. Since the efficiencies of phonon generation in electron and nuclear recoils are the same, the phonon channel provides direct calorimetry for both types of recoil. Super conducting elements sitting near phase transition provide extremely sensitive measurements of phonons.
EDELWEISS and CDMS are two experiments that measure phonons and ionization for discrimination in germanium crystals at millikelvin temperatures. The ionization yield for nuclear recoils is about four times lower than that for electron recoils. EDELWEISS-I has been operating in the Laboratoire Souterrain de Modane (LSM), which provides 4800 m.w.e. of rock shielding [186]. For temperature measurements, the detectors use neutron-transmutation-doped-germanium (NTD-Ge) thermometric sensors. The detectors also have aluminum electrodes for establishing the charge collecting field. See Figure 2.15 for a diagram of the detector. To select events in the center where the fields are more uniform, a cut is applied which requires that >75% of the charge signal be seen by the middle electrode. More details on the detector can be found in [187]. This detector achieved a rejection efficiency of >99.9% while maintaining a 90% acceptance for nuclear recoils down to an energy of 15 keV. An exposure of 62 kg-days was accumulated with three 320 g modules over 3 years and a cross section limit was published [186]. The collaboration is currently testing the twenty-eight 320 g Ge modules of their next generation detector called EDELWEISS-II. With a new cryostat, they hope to include up to 100 detectors and gain a two orders of magnitude improvement in sensitivity over EDELWEISS-I [188].
CDMS developed their first detectors in the 1990s, which ran in the Stanford Underground Facility. The later detectors, which have been running in the Soudan Underground Laboratory of Minnesota, have consistently given the best cross section upper limits. CDMS uses Si and Ge detectors which have >99.99% background discrimination efficiency above 10 keV in the detector’s bulk region. During the initial operation of their 160 g Ge detector, CDMS-I, many events were observed to have incomplete charge collection and, as a result, resembled nuclear recoils. These events were determined to be occurring near the surface and were minimized by adding new ionization contacts [189]. A limit was published in 2000 with data from Si and Ge detectors [190]. Then in 2004, the collaboration published the first results of CDMS-II running in the Soudan Underground Laboratory, which collected a net Ge exposure of 19.4 kg-day and set what was the lowest exclusion limit at the time [191]. Their newest limit was set from a 121.3 kg-day exposure which observed no events in nineteen 250 g Ge and eleven 100 g Si modules [14]. The data gave the same minimum WIMP-nucleon spin-independent cross section as that of XENON10 (which set the lowest upper-limit at time of publication) and provides the best sensitivity for masses above 42 GeV (see Figure 2.24.) They have proposed 27 kg Ge experiment called SuperCDMS which is significantly larger than CDMS-II [192].

CRESST-I consisted of several 262 g sapphire crystals, each with a tungsten superconducting phase transition (SPT) thermometer [193]. A first series of runs in 1999 revealed a much higher than expected background. The background did not appear to be due to radioactivity as it was time dependent and not Poissonian. The source of the background was eventually identified as cracks in the sapphire crystal. The cracks formed due to the tremendous amount of force applied at the small contacts. The problem was remediated by replacing the sapphire contacts with plastic ones. Using the crystal with the lowest energy threshold and a second one to eliminate coincidence events, 140 hours of data was collected to set a cross section limit [193].

CRESST-II at Gran Sasso National Laboratory takes a different approach by measuring scintillation and phonons in Ca$_2$WO$_4$ [194]. Ca$_2$WO$_4$ was chosen because of its relatively
high light yield and lack of degradation in light yield for events near the edge. The difference between the scintillation output of nuclear and electron recoils provides the discrimination. Similar to the measurement in CRESST-I, the phonons are measured with the tungsten SPT thermometers. The light is measured with a separate silicon wafer containing another SPT thermometer. To increase the collection efficiency of phonons created by photons, there are additional aluminum phonon collectors. To further increase the light collection, a 99% reflective polymeric multilayered foil surrounds the crystal. The group has taken 20.5 kg-days of data with two Ca$_2$WO$_4$ 300 g prototypes to set a WIMP-nucleon cross section limit [194].

2.3.5 Liquid Noble Element Detectors

Detectors using liquid neon, argon or xenon can be practically enlarged to the 1-ton scale unlike competing technologies such as cryogenic semiconductors. Furthermore, liquid noble elements are relatively inexpensive, easily purified, require relatively simple cryogenics, have high scintillation and charge yields, allow for background discrimination and can provide self-shielding. Also, liquid neon and xenon do not have problematic long-lived radioactive isotopes. On the other hand, argon has a large background from $^{39}$Ar which beta decays with an endpoint energy of 565 keV and half-life of 269 years and is generated by cosmic ray interactions via $^{40}$Ar(n, 2n)$^{39}$Ar. Its presence in commercial argon results in a background rate of $\sim$1 Bq/kg. Liquid argon detectors must rely on the highly efficient pulse shape discrimination to overcome this background. Xenon has $^{85}$Kr which beta decays with a half-life of 10.76 years but can be removed with distillation or chromatography.

Here, we describe the two basic forms of these detectors – single- and two-phase.

Single Phase Detectors

Single phase detectors measure only the scintillation light, which can provide background discrimination if the detector is sensitive to the photon arrival times. The scintillation pulse of nuclear recoils tend to have a larger fraction of the light at the head of the pulse
Table 2.1: Scintillation constants of liquid neon, argon and xenon [71, 22, 72].

<table>
<thead>
<tr>
<th>Element</th>
<th>$\tau_1$ (ns)</th>
<th>$\tau_2$ (µs)</th>
<th>$I_1$</th>
<th>$I_3$</th>
<th>ER</th>
<th>NR</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ne</td>
<td>18.2</td>
<td>14.9</td>
<td>0.10</td>
<td>0.66</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ar</td>
<td>4.0</td>
<td>1.59</td>
<td>0.3</td>
<td>3</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Xe</td>
<td>2.2</td>
<td>27</td>
<td>0.05</td>
<td>1.6</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Particles which promote faster recombination and hence shorten the timescale for photon emission from P2. A similar behaviour is expected for the high density tracks left by nuclear recoils, and this is confirmed by neutron calibrations. This difference in pulse times is the basis of the ZEPLIN I detectors described in this paper. More powerful discrimination can be achieved by using both the scintillation and ionisation processes in P2, extracting the ionisation to produce both a primary scintillation signal and a secondary proportional scintillation signal for each event [8, 9]. This provides the basis of two larger and more complex detectors, ZEPLIN II and ZEPLIN III, currently under construction as a continuation of this programme [10].

### 3. ZEPLIN I detector design

The principal features of the detector are shown in Fig. 1. The target consists of a total mass of liquid xenon, with a fiducial mass contained in a 15 cm diameter copper vessel viewed from the top through quartz windows by three 8 cm quartz-window ETL9265Q photomultipliers (PMTs). The vessel is lined internally with 5 mm thick diffuse PTFE reflector. The PMT connections pass through housings forming part of an outer copper vacuum chamber. The viewing windows are mounted on 4 cm long turrets, lined internally with Al reflector. The purpose of the turrets is to provide active shielding zones to reject low energy gammas (and also neutrons) from residual U/Th activity in the PMTs.

Higher energy gammas producing low energy deposits in the target by small angle Compton scattering are registered with an estimated 90% efficiency (from a Monte Carlo simulation) in the surrounding 30 cm thick liquid scintillator coincidence veto for an energy threshold 100 keV. Outside this is a 25 cm thick lead shield to attenuate external gammas. The liquid scintillator also serves as a neutron moderator/absorber, and a high purity passive gamma shield. For each interaction in the liquid xenon, scintillation light reaching the three PMTs produces photoelectron (pe) signals which are integrated and digitised at 1 ns intervals.

Initial runs used a DAQ system based on a GPIB-controlled LeCroy LC475 oscilloscope, later runs using a PCI-based Acquiris DC270 digitiser. For both DAQ systems the trigger was pro-

![Fig. 1. Arrangement of ZEPLIN I detector (centre), veto and shielding for runs 1 and 2. For run 3, a 10 cm layer of hydrocarbon neutron shielding was added over the top Pb shielding.](image)

as nuclear recoils tend to generate a greater portion of excimers in the singlet state which decays more quickly than the triplet states. Table 2.1 shows the scintillation parameters of liquid neon, argon and xenon. The pulse shape is usually quantified by the tail fall-time or by the fraction of pulse integral in the head of pulse. Pulse shape discrimination is very powerful in liquid neon and argon because of the great disparity between the singlet and triplet times. Our measurement of the pulse shape discrimination in liquid xenon is described in Section 5.3. See Section 3.4 for more details on the principle of pulse shape discrimination.

In the 1990s, the DAMA group engaged in basic R&D relevant to single phase liquid xenon detectors. Before taking dark matter data, the DAMA group measured two key char-

![Figure 2.16: The Zeplin-I detector surrounded by veto and shield [12].](image)
acteristics of liquid xenon: the scintillation efficiency of nuclear recoils relative to electron recoils and the pulse shape of the two types of recoils. In 1996, DAMA collected dark matter data with a three PMT detector containing 6.5 kg of $^{85}$Kr-free xenon enriched in $^{129}$Xe to 99.5%. The detector collected a 1763.2 kg-day exposure and set a cross section limit similar to that set by their NaI detector. Pulse shape discrimination was employed to remove some of the background [195].

A similar detector called ZEPLIN-I was built by the UK Dark Matter Group. It had a fiducial mass of 3.2 kg and collected scintillation with three PMTs. Figure 2.16 is a diagram of the detector, shield and liquid scintillator veto. The long volume of liquid xenon between each of the PMTs and the active volume allowed for the rejection of background events from the PMTs by demanding that the signal be present in multiple PMTs. The pulse shape parameter $\tau_m$ was the time interval between 5% and 90% of the maximum amplitude in the tail. Figure 2.17 shows the spectrum of $\tau_m$ for electron recoils of 7–10 keV. Figure 2.18 shows the $\tau_m$ spectrums from data containing both types of recoils at various energy ranges. Data of only nuclear recoils was not taken. The detector collected 293 kg-days of data in the 2300 m.w.e. deep UK Boulby mine and set a 90% confidence limit for spin-independent cross section with a minimum at $1.1 \times 10^{-6}$ pb [12].

Another single-phase liquid xenon experiment is XMASS, which hopes to construct a multi-ton spherical detector for detecting solar neutrinos, neutrinoless double beta decay ($^{136}$Xe $\rightarrow^{136}$Ba + 2$\nu^{-}$) and WIMPs [196]. It will rely on self-shielding to remove background. A large number of PMTs point toward the center which allows for the vertex reconstruction that is required to remove multiple scatters and edge events. The group has built a smaller cube-shaped 100 kg (3 kg fiducial volume) prototype with 54 PMTs (9 on each face) demonstrating self-shielding and achieving a light collection of 0.6 photoelectrons/keV [196]. The group has also built a distillation system that has reduced $^{85}$Kr levels in xenon from about 3 ppb to 3.3 ppt [197]. In 2007, XMASS received funding for a 800 kg (100 kg fiducial volume) detector [198, 199]. They have started construction for this detector which is expected to be completed within two to three years. This detector
vided by a threefold coincidence using CAEN N417 discriminators set at single-pe levels. A separate digitiser channel was used for the summed signals from the 10 veto PMTs.

The chamber is cooled by an IGC PFC330/C212 Polycold/C213 refrigerator, and Xe gas is progressively condensed into the chamber through a Messer-Griesheim/Oxisorb/C213 purifier. A stainless steel dump vessel stores the evaporated Xe in the event of refrigeration failure. Full practical and operational details of the complete system will be described in a separate paper. This paper discusses only those topics relevant to the estimation of dark matter signal limits.

4. Time constant and pulse selection

Prior to underground operation the system was characterised by tests in a surface laboratory. Using gamma sources 57Co and 137Cs located outside the chamber, the energy calibration was found to be in the range 1.5–2.5 pe (total) per keV for gammas from 57Co interacting in the lowest part of the vessel (position dependence being discussed in Section 5). This figure varied with cooldown cycle, possibly due to variations in xenon purity (see Section 7).

Since the dark matter limit depends entirely on events with energies below 10 keV, the corresponding pulses contained pe numbers \( n_{pe} < 15–25 \) and, in integrated form, have a stepped appearance due to variations in xenon purity (see Section 7). The value of \( t_m \) at the peak of the distribution has a physical associated with the characteristic parent time constants of the emission process. The width of the distribution depends on the mean number of photoelectrons available.

Fig. 2. Example of observed time constant distribution of underground data pulses (diamond points) from a gamma-ray source energy range 7–10 keV, showing also the fit using Eq. (2) (dashed line). Distributions of identical shape are observed in laboratory tests using gamma-ray sources. Nuclear recoil events would be observable as a second population of events on the left hand tail of the distribution.

will have 812 hexagonal face PMTs in a pentakis dodecahedron configuration and sit in a 10 meter high and 10 meter diameter water shield with several PMTs for muon rejection.

The DEAP/CLEAN collaboration is working on several detectors with liquid neon and liquid argon. They have measured the scintillation parameters for liquid neon [72] and argon [200] in small prototypes and have constructed DEAP-I, a 7 kg liquid argon module with two PMTs that is currently running at SNOLAB [201]. They plan to construct a 3600 kg liquid argon module in 2009 [202].

Two-Phase Liquid Noble Element Experiments

Two-phase liquid noble element time projection chamber detectors measure both scintillation and ionization to discriminate against background gammas. An excellent summary of two-phase detectors is given by [37]. Figure 2.19 shows the operation of the two-phase detector. A particle collides with the target and generates excitations and ionization. Immediate de-excitation and recombination releases the “primary” scintillation. With an electric field established by biased mesh electrodes, a portion of the electrons are prevented from recombining and are drifted away from the event site. The electrons are pulled out of the liquid into the gas where they excite the xenon gas, releasing a second burst of light known
and 60 keV. These are in reasonable agreement with [13] but lower than values obtained in a similar energy range by the Rome group of $f_{\text{Xe}} = 0.45 \pm 0.12$ using a neutron generator, and $f_{\text{Xe}} = 0.65 \pm 0.10$ using an Am–Be source [14]. The latter result from the Rome group relied on a Monte Carlo simulation fitted to the measured nuclear recoil spectrum and treating the quenching factor as a free parameter. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.

The above measurements provide the quenching factor at $\sim 40$ keV recoil energy ($\sim 8$ keV electron-equivalent) but do not extend to our experimental threshold of 10 keV recoil energy (2 keV electron-equivalent). Theoretical discussions [15] do not suggest a definite behaviour of $f_{\text{Xe}}$. The same technique has been applied to the ZEPLIN-I tagged Am–Be data used in the top plot of Fig. 3 between 15 and 70 keV electron-equivalent energies. This gives a value of $f_{\text{Xe}} = 0.25 \pm 0.05$ averaged over a recoil energy range 60–280 keV.
Figure 2.19: Diagram of a two-phase detector with two PMT arrays. A particle interaction emits light and electrons. The electrons are pulled out of the liquid where they excite the gas and emit a burst of light. The light is usually detected by PMTs.
as the proportional or secondary scintillation. The fraction of electrons that are pulled out of the liquid depends on the electric field (see Section 3.5.3). Typically several tens of photons are generated per electron (see Section 3.5.4 for more details). This powerful technique has the ability to sense single electrons (an alternative method of directly measuring charge avalanches is described in Section 5.5). Both the primary and proportional light are measured by PMTs or similar devices. Typically, there is an array of PMTs in the liquid below the active volume and another array above in the gas. The PMTs in the liquid have a much greater collection efficiency for primary light than those in the gas because most of the photons stay in the liquid due to a mismatch in the refractive indices [42].

A multiple-PMT array in the gas allows for x-y position reconstruction. Typically, the primary and proportional scintillation signals are referred to as “S1” and “S2”, respectively, and the size of these signals as S1 and S2, respectively. The basis of charge-based discrimination (as with pulse shape discrimination) lies in the linear energy transfer (LET) – the fraction of charge recombination is positively correlated with the LET. This manifests as a lower ionization to scintillation ratio for nuclear recoils: 

\[(S2/S1)_{ER} > (S2/S1)_{NR},\]

where ER and NR stand for electron and nuclear recoils, respectively. In two-phase liquid argon detectors, both recombination and pulse shape discrimination are employed, with most of the discrimination power deriving from the latter principle. (Section 5.3 demonstrates the use of both principles simultaneously in LXe detectors.)

There are two two-phase experiments using argon. WARP has been running a small 2.3 liter (1.87 liter sensitive) module in the Gran Sasso National Laboratory. The detector has seven 2" PMTs in the gas and three mesh electrodes. The drift volume is surrounded by a high performance diffuse reflector coated with tetra-phenyl-butadiene (TPB), a wavelength shifter for converting the 128 nm photons to about 440 nm. Increasing the wavelength boosts the primary light signal by raising the detector reflectance and PMT quantum efficiency. Figure 2.20 shows \(\log(S2/S1)\) versus recoil energy with the single-hit nuclear recoils highlighted. Cross section limits were published based on an 96.5 kg-day exposure [203]. Their results indicated that the two discrimination methods are independent, which is nec-
necessary for the discrimination efficiencies to compound. They are currently building a 140 kg module surrounded by an argon anti-coincidence shield. To further reduce background they have sought argon with lower concentrations of $^{39}$Ar. Instead of pursuing centrifugation, which is likely to be prohibitively expensive, they have looked for underground sources that have been shielded from cosmic rays. One such source contains an $^{39}$Ar concentration that is 5% that of normal argon [204].

The other two-phase liquid argon detector is the 1 ton detector by ArDM which is currently under construction. The ArDM detector is similar to WARP except that the top PMT array is replaced with a GEM charge readout system (see Section 5.5.2) to reduce
Figure 2.21: Plot of log($S_2/S_1$) versus prompt/total between 40 and 60 keV, of the AmBe (gamma and neutron) and WIMP search data from WARP. The acceptance window is represented by the red box. Eight events were observed in the acceptance window but only five events remain in the single-hit window.

background [205, 206]. They have measured a gain of up to $10^4$ with a two-stage GEM structure in pure argon at 87 K.

There have been several two-phase liquid xenon dark matter detectors. XENON10 at Gran Sasso National Laboratory was the first such detector to publish a limit – this detector and its results are be discussed in Chapter 6. The ZEPLIN collaboration has published cross section limits for two two-phase liquid xenon detectors. In 2007, the collaboration published spin-independent [207] and spin-dependent [62] cross section limits for the 31 kg (7 kg fiducial volume) ZEPLIN-II detector at the Boulby Mine. Their next detector, ZEPLIN-III published limits about a year later. This 6.5 kg fiducial volume detector achieved better results than ZEPLIN-II as it had a PMT array in the liquid (as opposed in the gas) which gave a significantly higher light collection and lower energy threshold [208]. Their first run of 847 kg-days exposure set a spin-independent cross section limit with a minimum of $7.7 \times 10^{-8} \text{pb}$ at the WIMP mass of 55 GeV [15]. This data was also used to set a spin-dependent WIMP-neutron cross section limit with minimum of $1.8 \times 10^{-2} \text{pb}$ at 55 GeV [16]. Both limits improved upon the results of XENON10 for masses above $\sim 100$ GeV. Liquid xenon detectors are unable to set competitive WIMP-proton limits because of the
lack of isotopes with unpaired protons. There are currently two 100 kg scale detectors that are likely to publish new limits within a year. The XENON collaboration is working on XENON100 which has a total mass of 170 kg and 65 kg in fiducial volume and is running in the cryostat and shield that previously housed XENON10 [209]. The LUX collaboration is working on 300 kg module with 100 kg fiducial volume which will run in the Homestake mine [210].

2.3.6 Gaseous Detectors

The DRIFT collaboration has been developing the low pressure Negative Ion Time Projection Chamber (NITPC) for WIMP detection. One goal of these detectors is to deduce the directionality of the WIMP “wind” by reconstructing the ion tracks. The detectors employ large time projection chambers with multi-wire readouts. The tracks are reconstructed by signal size and arrival times of the ionization. For this approach, gas is needed to generate ionization clouds large enough for their shapes to be probed. The target medium is $\sim 40$ torr carbon disulfide ($\text{CS}_2$), which is a electronegative gas that replaces electrons as the charge carrier and consequently lowers the loss of track information due to diffusion. Good discrimination on the order of $10^6$ is achieved between nuclear and electron recoils as they have drastically different track lengths. Figure 2.22 shows the wire signals for alpha and neutron events. In this figure, the number of ion pairs (NIPs) collected by each wire and track lengths (RDs) are given for each event, showing that the ionization density (total NIP/RD) is larger for the nuclear recoil event. The sulphur nucleus has a small spin-independent cross section and thus needs to be compensated by increasing the total target mass. The collaboration has finished constructing and testing DRIFT-I, a 1 m$^3$ module [211], and is working on DRIFT-II which runs with a higher drift field to decrease the effects of diffusion [212].
Figure 2.22: DRIFT-I wire signals of alpha (top) and neutron (below) events. The number of ion pairs (NIPs) collected for is indicated for each wire and is proportional to the energy collected by the wire. The “RD” is the 2D track length. The ionization density (total NIP/RD) is larger for nuclear recoil event.
2.3.7 Superheated Liquid Bubble Detectors

The invention of the bubble chamber in the 1950s by Donald Glaser lead to the discovery of many particles [213]. The application of this concept for cold dark matter detection was first proposed in the 1990s [214] and has been implemented in the search for dark matter. These detectors employ superheated fluids in which bubbles form upon energy deposition of sufficient concentration. Several materials have been used which allow nuclear recoils (and other interactions of similar LET) to form bubbles. In additional to having very high discrimination efficiencies (>10^9), these detectors are cheap, scalable, and require fairly simple equipment. One of the disadvantages is the presence of alpha-emitting contaminants and the spontaneous formation of bubbles by impurities and surface defects. The bubbles are usually detected with microphones and their positions can be reconstructed with CCD cameras. There are two basic approaches. The superheated droplets detectors (SDDs)
have micro-droplets (∼10 µm) suspended in a matrix. The other approach employs bulk superheated fluids. The former technique helps to avoid problems with the spontaneous formation of bubbles at the edges. This technique, however, is limited by which targets can be used because of the need to avoid de-emulsification. These experiments are limited to targets such as C₂ClF₅, which are good for spin-dependent searches due to the presence of fluorine but are not good for spin-independent searches. CF₃I and CF₃Br are good for both searches but cannot be used well in SSDs. They have to be used in bulk superheated fluid detectors which generally have problems with surface defects but allow for easily interchanging of target material.

PICASSO and SIMPLE are two SDD experiments. SIMPLE has published a spin-dependent cross section limit comparable to other detectors despite its relatively low 0.42 kg-day exposure with C₂ClF₅ [215]. SIMPLE has also explored the use of CF₃I but has yet to achieve long term stability. PICASSO has published spin-dependent cross section limits with a 2 kg-day exposure using C₄F₁₀ [65] and is currently working on thirty-five 4.5 liter modules each with a mass of 2.6 kg [216].

The Chicagoland Observatory for Underground Particle Physics (COUPP) experiment uses a bulk superheated fluid. To lessen the problem of surface defects, smooth quartz vials are used with a lining of buffer fluid. The wetting of the surface can be improved by slow vapor deposition, and the trapped bubbles in the surface defects can be outgassed. With a set of targets of different spin-dependent and spin-independent sensitivities (for example, C₃F₈ or C₄F₁₀, which have higher spin-dependent couplings than CF₃I), they can make more precise measurements of the coupling constants [217]. Figure 2.23 shows the tracks of neutrons and muon events in a COUPP prototype. COUPP has collected a 250 kg-day exposure with 1.5 kg CF₃I module running underground in Fermilab which was able to improve spin-dependent WIMP-proton cross section limits even with little attention paid to alpha-recoil backgrounds from U/Th in quartz and radon from air, rubber seals, and welds [66]. This detector demonstrated a background leakage of ∼10¹⁰ at ∼10 keV recoil energy which is unmatched by any other WIMP detector; for comparison, CDMS has a
Figure 2.24: 90% confidence upper limit to the WIMP-nucleon elastic scattering cross sections set by the top three experiments: XENON10 [13], CDMS-II [14] and ZEPLIN-III [15]. Figure taken from [15].

rejection factor of $10^{-4}$–$10^{-5}$. COUPP is currently constructing a $\sim$50 kg detector.

Figure 2.24 shows the spin-independent cross section upper limit set by XENON, CDMS and ZEPLIN-III. Figure 2.25 shows the upper limits on pure WIMP-neutron and WIMP-proton cross sections set by many experiments. Two-phase liquid xenon detectors (XENON and LUX) are in a position to continue leading the field in setting spin-independent and WIMP-neutron limits as they do not have any technical roadblocks against further expansion. Furthermore, as they scaled up, they will also be taking increasing advantage of self-shielding. Unlike single-phase liquid xenon detectors, two-phase detectors can efficiently remove multiple scatter neutron events. Argon experiments must find argon with low intrinsic background to be competitive. CDMS has consistently generated top results but the expansion to SuperCDMS would be expensive (> $30 million). Gaseous detectors like DRIFT are unlikely to set new limits due to their low mass. The droplet detectors have enormous potential due to their nearly infinite background rejection efficiency, low instru-
especially
Our xenon was depleted from high-mass isotopes, es-
odd-neutron isotopes reflects their relative abundance.

The exposure of the
and nuclear recoil energies utilises the varying quenching
energy, is applied separately. Conversion between visible
the restricted acceptance). An energy-dependent detec-
tion efficiency, which reaches unity near 14 keV recoil

ments place the world's most stringent limit on
σ
these limits unfavourably:
this
with characteristic velocity

ter halo with truncated Maxwellian velocity distribution
scribed Ref. [17]. We assumed an isothermal dark mat-
ments.

The SD cross-section limits, with no assumption on the
SD cross-section set by many experiments: ZEPLIN-I, ZEPLIN-II, XENON10, 
Figure 2.25: Upper limits to pure WIMP-neutron (top) and WIMP-proton (bottom) spin-
proton spin-depended cross-sections. In addition to ZEPLIN-III XE, both
in potentials, we show other xenon experiments in
black: ZEPLIN-I [22], ZEPLIN-II [23] and XENON10 (Bonn
A) [18, 19]. Additional curves are CDMS-II [24], COUPP [25],
EDELWEISS [26], KIMS [27], NAIAD, PICASSO [29] and SIMPLE [30]. The pure-proton indirect limit from Super-Kamiokande is also shown [31]. The DAMA evidence region
interpreted as a nuclear recoil signal in a standard halo [7] is
indicated in green. The hatched area is the tip of the 95% probability region for neutralinos in the Constrained Minimal Supersymmetric Standard Model (CMSSM) [17]. Figure taken from

FIG. 1: Upper limits on pure WIMP-neutron and WIMP-proton SD cross-sections. In
addition to ZEPLIN-III XE, both in potentials, we show other xenon experiments in
black: ZEPLIN-I [22], ZEPLIN-II [23] and XENON10 (Bonn A) [18, 19]. Additional curves are CDMS-II [24], COUPP [25], EDELWEISS [26], KIMS [27], NAIAD, PICASSO [29] and SIMPLE [30]. The pure-proton indirect limit from Super-Kamiokande is also shown [31]. The DAMA evidence region
interpreted as a nuclear recoil signal in a standard halo [7] is indicated in green. The hatched area is the tip of the 95% probability region for neutralinos in the Constrained Minimal Supersymmetric Standard Model (CMSSM) [17]. Figure taken from

that is derived from a
fiducial 'geometrical' ex-
subsuming all energy-
deficiencies (and
energy-dependent detector-
near 14 keV recoil
ersion between visible
varying quenching

The exposure of the
relative abundance.

e Xe used in 0
νββ
=600 km/s, Earth velocity

by orders of magnitude, virtually elim-
ning the sensitivity to WIMP-proton scattering (this

The allowed region of
parameter space can be
defined by
completing the square:

FIG. 1: Upper limits on pure WIMP-neutron and WIMP-
proton spin-depended cross-sections. In addition to ZEPLIN-III XE, both
in potentials, we show other xenon experiments in
black: ZEPLIN-I [22], ZEPLIN-II [23] and XENON10 (Bonn
A) [18, 19]. Additional curves are CDMS-II [24], COUPP [25],
EDELWEISS [26], KIMS [27], NAIAD, PICASSO [29] and SIMPLE [30]. The pure-proton indirect limit from Super-Kamiokande is also shown [31]. The DAMA evidence region
interpreted as a nuclear recoil signal in a standard halo [7] is
indicated in green. The hatched area is the tip of the 95% probability region for neutralinos in the Constrained Minimal Supersymmetric Standard Model (CMSSM) [17]. Figure taken from

...
mental complexity and very low cost. With the elimination of alpha-emitting contaminants, COUPP may see an improvement in sensitivity of up to six orders of magnitude, excluding much SUSY space.

2.4 High Energy Collider Searches

The discovery of the neutralino could be one of the major discoveries at the Large Hadron Collider (LHC) at CERN [218, 219]. The 14 TeV proton collider is expected to provide enough energy for such particles if they exist. These particles are likely to escape undetected and be identified with a missing mass signature. The future International Linear Collider (ILC) will also be able to search for neutralinos via 1 TeV electron-positron collisions [220]. Due to a reported price tag of about $7 billion [221], this accelerator may never be constructed. It must be noted that even if the neutralino is discovered in a collider, indirect and direct detection experiments will still be needed to show the existence of a neutralino halo.
Chapter 3

Xenon

3.1 Introduction

Xenon is a colorless and odorless noble gas of atomic number 54 that exists in trace amounts in the atmosphere. This noble gas was discovered on July 12, 1898 by Sir William Ramsay and Morris Travers, who named it after the Greek word ξένον, meaning foreign or strange. Since its discovery, many practical uses of xenon have been found in medicine, lighting, propulsion and particle detection. Major research on the application of liquid xenon in multi-wire ionization chambers were performed in the 1960s. Throughout the next two decades, measurements were done on key values such as the Fano factors, scintillation yields, and decays times. A detailed history of xenon detectors can be found in [222]. In this section, we give a brief overview of some of the applications of xenon, its general physical and chemical properties, and specific properties related to its application in particle detectors.

3.2 Applications of Xenon

Xenon has been used since the late 1970s as a propellant in ion thrusters of spacecrafts because of its low ionization potential per atomic weight [223]. In these devices, xenon ions are first created by bombarding xenon with a stream of electrons and then accelerated by an
Figure 3.1: Diagram of an ion propulsion system. Xenon is typically used as the propellant because of its low ionization potential per atomic weight.

electric field. The electrons are kept in the ionization region by a magnetic field (See Figure 3.1). These thrusters require $\sim 10$ kW of power and generate a thrust of several hundred mN.

Xenon gas is an excellent light source. It has been used in flash lamps because of its fast response and also in movie projector lamps because their emission spectrum resembles natural light more than that of typical light bulbs. Xenon-neon mixtures are used in plasma television screens.

Xenon serves multiple purposes in medicine. The anesthetic effects of xenon on mice were first published in 1946 [224]. Soon afterward, its anesthetic effects were tested in humans [225] but widespread use has only recently been seriously considered [226]. It has many properties of an ideal anesthetic: high insolubility in human tissue enabling faster induction and emergence than other inhaled anesthetics, low toxicity, and absence of teratogenicity. Furthermore, the use of xenon has been motivated by environmental concerns as it does not destroy the ozone layer like common anesthetics. Although expensive, the use of xenon
Figure 3.2: Position emission tomography. Liquid xenon detectors on the periphery detect the gammas emitted from position-electron annihilations.

can be cost effective if recycled in a closed environment.

Xenon is also used in medical imaging both as a source and as a detector medium. A key feature of xenon is its rapid clearance from the body, which importantly limits the dosage to the patient. One such diagnostic inhalant is $^{133}$Xe which has a half-life of 5.25 days and is produced in the fission of $^{235}$U. Only the 80.9 keV gamma ($\Gamma = 38.0\%$) is used for imaging. $^{133}$Xe has been used to successfully monitor blood flow [227]. $^{129}$Xe has been used in magnetic resonance imaging (MRI). By hyperpolarizing $^{129}$Xe, the nuclear magnetic resonance signal is greatly enhanced. Hyperpolarization of $^{129}$Xe to four or five orders of magnitude above the Boltzmann polarization can be achieved by having the xenon undergo spin exchange with evaporated rubidium that has been polarized with circularly polarized light [228]. In the gas phase, the $^{129}$Xe has a relaxation time of several hours [229] while in blood, the polarization lasts for only several seconds [230]. Liquid xenon has been considered as a target in positron emission tomography (PET), which is a medical imaging technique whereby the patient is injected with a positron-emitting tracer isotope. The two
Table 3.1: General properties of noble elements.

<table>
<thead>
<tr>
<th>Property</th>
<th>He</th>
<th>Ne</th>
<th>Ar</th>
<th>Kr</th>
<th>Xe</th>
</tr>
</thead>
<tbody>
<tr>
<td>Symbol</td>
<td>He</td>
<td>Ne</td>
<td>Ar</td>
<td>Kr</td>
<td>Xe</td>
</tr>
<tr>
<td>Atomic Number</td>
<td>2</td>
<td>10</td>
<td>18</td>
<td>36</td>
<td>54</td>
</tr>
<tr>
<td>Atomic Radius (pm)</td>
<td>31</td>
<td>38</td>
<td>71</td>
<td>88</td>
<td>108</td>
</tr>
<tr>
<td>Gas Density at 1 atm, 298 K (g/l)</td>
<td>0.16</td>
<td>0.82</td>
<td>1.63</td>
<td>3.43</td>
<td>5.40</td>
</tr>
<tr>
<td>Liquid Density at 1 atm, $T_{boiling}$ (g/cm$^3$)</td>
<td>0.12</td>
<td>1.21</td>
<td>1.40</td>
<td>2.41</td>
<td>2.94</td>
</tr>
<tr>
<td>Boiling Temperature at 1 atm (K)</td>
<td>4.22</td>
<td>27.07</td>
<td>87.30</td>
<td>119.93</td>
<td>165.03</td>
</tr>
<tr>
<td>Melting Temperature at 1 atm (K)</td>
<td>-</td>
<td>24.56</td>
<td>83.80</td>
<td>115.79</td>
<td>161.4</td>
</tr>
<tr>
<td>Critical Temperature (K)</td>
<td>5.19</td>
<td>44.4</td>
<td>150.8</td>
<td>209.4</td>
<td>289.7</td>
</tr>
<tr>
<td>Critical Pressure (atm)</td>
<td>2.24</td>
<td>27.2</td>
<td>48.1</td>
<td>54.3</td>
<td>57.6</td>
</tr>
<tr>
<td>Critical Density (g/cm$^3$)</td>
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<td>0.482</td>
<td>0.544</td>
<td>0.91</td>
<td>1.10</td>
</tr>
<tr>
<td>Triple Point Temperature (K)</td>
<td>N/A</td>
<td>24.56</td>
<td>83.81</td>
<td>115.78</td>
<td>161.41</td>
</tr>
<tr>
<td>Triple Point Pressure (atm)</td>
<td>N/A</td>
<td>0.424</td>
<td>0.68</td>
<td>0.722</td>
<td>0.805</td>
</tr>
<tr>
<td>Dielectric Constant of Liquid</td>
<td>1.049</td>
<td>1.188</td>
<td>1.51</td>
<td>1.66</td>
<td>1.95</td>
</tr>
</tbody>
</table>

collinear 511 keV gamma rays from positron-electron annihilations are detected by a ring of modules containing PMTs and electrodes in LXe (as seen in Figure 3.2) [231, 232]. The annihilation positions are reconstructed by time-of-flight information.

### 3.3 General Properties

Xenon is found in the atmosphere at concentration of 0.09 ppm and can be extracted by fractional distillation, requiring 220 watt-hours of energy per gas liter. The melting and boiling points of xenon at 1 atm are 161.4 K and 165.0 K, respectively, and as for all noble elements (except helium, which does not freeze at 1 atm), the temperatures are close. Figure 3.3 shows the phase diagram of xenon and Figure 3.4 shows the vapor pressure as a function of temperature. Of all the noble elements, xenon in liquid form has the second highest density at about 3.0 g cm$^{-3}$ (radon has a density of 4.4 g cm$^{-3}$). Figure 3.5 is a plot of the liquid xenon density versus temperature. Solid xenon has a density of 3.64 g cm$^{-3}$.

As with all noble elements, xenon undergoes very few chemical reactions. Some examples of synthetic xenon compounds are PtF$_6$Xe, XeF$_2$, XeF$_4$, and XeF$_6$. Some basic properties of xenon and other noble elements are given in Table 3.1.

There are nine stable naturally occurring isotopes: $^{124}$Xe, $^{126}$Xe, $^{128}$Xe, $^{129}$Xe, $^{130}$Xe,
Figure 4.1: Phase diagram of Xe.

Figure 3.3: Phase diagram of xenon.

$^{131}$Xe, $^{132}$Xe, $^{134}$Xe, and $^{136}$Xe. Only tin has more stable isotopes. The properties of these isotopes are summarized in Table 3.2. $^{124}$Xe, $^{134}$Xe and $^{136}$Xe are expected to double beta decay but have not been observed to do so. The Enriched Xenon Observatory (EXO) aims to detect the neutrino-less double beta decay of $^{136}$Xe, which would prove that neutrinos are massive Majorana particles [233]. Over 40 isotopes have been created and studied. $^{131m}$Xe, $^{133}$Xe, $^{133m}$Xe, and $^{135}$Xe result from fission of $^{235}$U and $^{239}$Pu making them indicators of nuclear explosions. $^{135}$Xe, with its large cross section to thermal neutrons ($2.65 \times 10^6$ barns), is responsible for the “xenon poisoning” seen in reactors whereby $^{135}$Xe captures the neutrons (and transforms into stable $^{136}$Xe) that would otherwise initiate fission. This isotope is generated by the decay of $^{135}$I, a common fission product. Under normal operations, the reaction rate is set such that $^{135}$Xe is burned and maintained at an equilibrium
Figure 3.4: The vapor pressure of xenon as a function of temperature.
concentration. A failure to recognize the presence of $^{135}$Xe caused the Chernobyl disaster [234].

Although, xenon has no long-lived isotopes, it does contain $^{85}$Kr, which is a beta emitter. Fortunately, $^{85}$Kr can be removed to negligible levels by distillation [197] or chromatography [235].

### 3.4 Basic Particle Interactions with Xenon

An important characteristic of the target material is its efficiency in transforming radiation energy into measurable quanta (e.g., phonons, photons, and electrons). The type of interaction depends on the particle’s charge, mass, and momentum. Light charged particles such as electrons and alphas interact electromagnetically producing tracks of atomic excitations and ionizations. Charged particles can also generate bremsstrahlung and/or Cerenkov radiation but neither of these processes occur in our energy range of interest (<100 keV).
Chargeless particles like gammas and neutrons do not directly generate bulk excitation and ionization directly but do so by generating recoiling electrons or nuclei, respectively. A recoiling heavy nucleus not only generates atomic excitations and ionizations but also atomic motion. This results in an apparent suppression (or “quenching”) of scintillation and ionization with respect to electron recoils of the same energy. It is important to point out that discrimination in two-phase xenon detectors do not rely on this suppression as both the scintillation and ionization channels are equally suppressed. Both photons and electrons can travel long distances in xenon that is free of impurities.

This section describes the basic interactions of various particles with xenon, the processes concerning the intermediate products and the transport properties of the induced ionization and scintillation.

### Gamma Ray Interactions

Gamma rays eject an electron through photoelectric absorption, Compton scattering or pair production (if the energy of the gamma exceeds twice the rest mass of the electron). The cross section of these three processes are proportional to $Z^4/E_\gamma^{7/2}$, $Z/E_\gamma$, and $Z^2 \ln(2E_\gamma)$, respectively. Photoelectric absorption results in the complete absorption of a photon, generating an electron with energy equal to that of the photon minus the ionization energy. In Compton scattering, only a fraction of the photon energy is spent in ejecting the electron.
Figure 3.6: The mass attenuation coefficient (top) for photons on xenon and the total attenuation length (bottom) of gammas in liquid xenon as a function of energy [18].
and that energy is related to the scattering angle by

\[ \frac{1}{E_f} - \frac{1}{E_0} = \frac{1}{m_e c^2} (1 - \cos \theta). \]  \hspace{1cm} (3.1)

Pair production is only possible for gammas with energy \( E_\gamma \geq 2m_e = 1.02 \text{ MeV} \), the rest mass of two electrons. Photoelectric absorption is the dominant interaction of lower energy photons. Figure 3.6 shows the mass attenuation coefficient for the three different types of gamma interactions and the total attenuation length of gammas in liquid xenon as a function of energy. The structures at \( \sim 5 \text{ keV} \) and \( \sim 30 \text{ keV} \) represent the L-shell and K-shell absorption edges. The abrupt drops in the mean free path at edges are due to the energy of the gamma being sufficient to eject electrons of an additional shell. The vacancy is filled with an other shell electron which releases an x-ray or Auger electron. These plots show the effectiveness of LXe for self-shielding. The attenuation length of 1 MeV gammas is \( \sim 6 \text{ cm} \) in LXe.

**Neutron Interactions**

Neutrons interact primarily with nuclei since, as chargeless particles, they can not interact with the surrounding electrons. Neutrons can either scatter off of or be absorbed by nuclei. In elastic scattering, the total kinetic energy remains the same. A neutron recoiling off a nucleus of atomic weight \( A \) will lose an average energy of \( 2EA/(A + 1)^2 \). This expression explains the effectiveness of hydrogen-rich materials (such as water and plastic) as neutron moderators. In the case of inelastic scattering, some of the energy can be channeled into exciting the nucleus. The excited nucleus eventually relaxes and radiates. A neutron can also be absorbed by nucleus. In general, the addition of a neutron can lead to a rearrangement of the nuclear structure and an eventual emission of gammas, neutrons, protons, deuterons or alphas. Fission fragmentation can also occur. Collisions yielding protons, deuterons, multiple neutrons and other larger fragments are negligible until \( E > 1–10 \text{ MeV} \). Some nuclear excited states are very brief such as \(^{125m}\text{Xe} \), which decays with a half-life of 57 s and emits a 253 keV gamma. \(^{125m}\text{Xe} \) can be produced by absorption of thermal neutrons, \((^{124}\text{Xe}(n,\gamma)^{125m}\text{Xe})\). \(^{131m}\text{Xe} \), which is produced by inelastic scattering of \(^{131}\text{Xe} \),
has a half-life of 11.8 days and emits 163.9 keV gammas upon relaxation. Its long half-life makes it a useful source for uniformly probing all regions of the detector [236].

**Recoiling Particles in Liquid Xenon**

In this section, we give a general overview of the interactions of recoiling particles in liquid xenon and the interactions of resulting products.

Energetic particles slow down as they deposit energy into the detector medium. They can lose energy electronically through the creation of atomic excitations and ionization or by generating atomic motion (phonons). Alphas and electrons (from gamma interactions and beta decays) lose energy almost entirely through the former process whereas a recoiling xenon atom (from a neutron or WIMP collision) will lose energy primarily through the latter.

The electronic energy deposition creates scintillation by the following processes:

1) De-excitation

\[
\text{Xe}^* + \text{Xe} \rightarrow \text{Xe}_2^* \\
\text{Xe}_2^* \rightarrow 2\text{Xe} + h\nu
\]

2) Recombination

\[
\text{Xe}^+ + \text{Xe} \rightarrow \text{Xe}_2^+ \\
\text{Xe}_2^+ + e^- \rightarrow \text{Xe}^{**} + \text{Xe} \\
\text{Xe}^{**} \rightarrow \text{Xe}^* + \text{heat} \\
\text{Xe}^* + \text{Xe} \rightarrow \text{Xe}_2^* \\
\text{Xe}_2^* \rightarrow 2\text{Xe} + h\nu
\]

The scintillation of all noble elements is believed to follow this mechanism. Importantly, this scintillation mechanism results in the generation of photons that do not ionize the xenon. In the presence of an electric field, some of the electrons are prevented from recombining which results in less scintillation. The recombination scintillation mechanism results in an
FIG. 3: Energy dependence of nuclear recoil ionization yield at different drift fields.

FIG. 4: Top: Field dependence of scintillation and ionization yield in liquid xenon for 122 keV electron recoils (ER), 56 keV nuclear recoils (NR) and alphas. Bottom: Ionization yields scaled by their 4.5 kV/cm values.

Electrons, alphas and Xe atoms propagate through liquid xenon in different ways manifesting as different charge and light yields. The top graph of Figure 3.7 shows the light and charge yield as a function of field for 122 keV gammas, 56 keV nuclear recoils, 5.3 MeV alphas relative to that of the maximum value as a function of electric field. Below – same iodine data (light yield) and a 122 keV electron recoils. Measurements were performed by the Case and Columbia groups [19].

Recombination is primarily a function of electric field and ionization density, with stronger recombination at low fields and in denser tracks. Ionization density along a track corresponds roughly to electronic stopping power, plotted in Fig. 5 for alphas, electrons, and Xe nuclei in LXe, as given by ASTAR, ESTAR, and SRIM [21], respectively.
FIG. 1: Columbia detector response to AmBe neutron (top) and 137Cs gamma sources (bottom), at 2 kV/cm drift field.

FIG. 2: Case detector response to 252Cf neutron (top) and 133Ba gamma sources (bottom) at 1.0 kV/cm drift field.

FIG. 3: Energy dependence of nuclear recoil ionization yield at different drift fields.

FIG. 4: Top: Field dependence of scintillation and ionization yield in liquid xenon for 122 keV electron recoils (ER), 56 keV nuclear recoils (NR) and alphas. Bottom: Ionization yields scaled by their 4.5 kV/cm values.

Recombination is primarily a function of electric field and ionization density, with stronger recombination at low fields and in denser tracks. Ionization density along a track corresponds roughly to electronic stopping power, plotted in Fig. 5 for alphas, electrons, and Xe nuclei in LXe, as given by ASTAR, ESTAR, and SRIM [21], respectively. Also shown is the total energy lost to electronic excitation per unit path length for Xe recoils, which differs from the electronic stopping power in that it includes energy lost via electronic stopping of secondary recoils. The decrease in Xe recoil stopping power at lower energies is responsible for the rise in ionization yield (see Figure 3.8, which shows measurements by the Case group and our collaborators at Columbia [19]). The stopping power of Xe nuclei at 56 keV (including the daughter recoils) is higher than that of 5.3 MeV alphas. Presumably this would translate
Figure 3.9: The electronic stopping power of electron, alpha and Xe atom recoils [18, 20, 21]. The Hitachi result includes the energy lost via stopping of secondary recoils.

into a higher charge density and thus a lower ionization yield for the nuclear recoils. Instead, we observe that the nuclear recoils give the higher ionization yield. This suggests that other aspects such as the track structure also play a role. Alphas are presumed to have straight cylindrical tracks with a dense core and a “penumbra” of delta rays [237]. A nuclear recoil track Monte Carlo was written using cross sections and stopping powers in [238] and is described in [60]. Simulations show that the nuclear recoil tracks have many branches (see Figure 3.10(a)) with much of the energy carried away by these branches, which presumably end in a very sparse track since the stopping power falls with energy. The electron recoil track (see Figure 3.10(b)), in comparison, does not have the branching structure of Xe recoils and are larger with lower ionization density, hence the relatively higher yield and field dependence. Electron recoil tracks were created with PENELOPE, a simulation package for propagating electrons and photons [239].

The scintillation pulses in liquid noble element detectors have been observed to contain
Figure 3.10: Tracks of 10 keV nuclear recoil (top) and 4 keV electron recoil (bottom) in liquid xenon. The ions are indicated by dots.
values. Kubota et al.14 and Carvalho and Klein also reported smaller values under \( \alpha \)-particle excitation.

The reason for this disagreement is not well understood at present. A possible reason is a difference in liquid temperature and impurities. The energy of excited states in argon is relatively large and those states are likely deexcited by collisions with impurities. According to our experience, nonpurified argon showed a shorter lifetime for \( ^3\Sigma_\pi^+ \).

Also, it has been reported that a few tens ppm of impurities such as nitrogen, oxygen, or carbon monoxide can lead to a substantial reduction in intensity of luminescence originating from low excited states.\(^{18}\)

The intensity ratios \( I_S/I_T \) of the singlet states \( \left( ^2\Sigma \right) \) to the triplet states \( \left( ^3\Sigma \right) \) are found to be 0.3, 1.3, and 3 for electron, \( \alpha \)-particle, and fission-fragment excitation, respectively, and this result shows an enhancement of \( ^1\Sigma \) formation in higher deposited energy density. This result shows the same trend as results obtained by Kubota et al.\(^{13,14}\) under electron and \( \alpha \)-particle excitation for condensed argon and xenon.

B. Liquid xenon

The decay curves obtained under \( \alpha \)-particle and fission-fragment excitation in liquid xenon showed double exponential decay forms. Typical decay curves obtained for liquid xenon are shown in Fig. 4. The lifetimes obtained for \( ^3\Sigma_\pi^+ \) under \( \alpha \)-particle and fission-fragment excitation are 4.3 \( \pm 0.6 \) and 4.3 \( \pm 0.5 \) nsec, respectively, and those for \( ^3\Sigma_g^+ \) are 22\( \pm 1.5 \) and 21\( \pm 2 \) nsec, respectively, and agreed within experimental errors. The values obtained under \( \alpha \)-particle excitation agreed with those reported by Kubota et al.\(^{11}\) The results are listed in Table 1.

The time dependence of luminescence from liquid xenon excited by electrons is shown in Figs. 4 and 5, and is quite different from those observed under \( \alpha \)-particle and fission-fragment excitation. The curve shows a relatively slow rise and does not follow an exponential form. Kubota et al.\(^{19}\) reported that this nonexponential component disappears when an electric field is applied. Therefore, it is clear that the nonexponential component is due to a contribution from the recombination.

The ratios \( I_S/I_T \) are found to be 0.45\( \pm 0.7 \) and 1.6\( \pm 0.2 \) under \( \alpha \)-particle and fission-fragment excitation, respectively, showing an enhancement of \( ^1\Sigma^+ \) formation with higher deposited energy density.

Figure 3.11: The scintillation pulse shape under different particle excitations at zero electric field as reported by Hitachi\(^{22}\).
KUBOTA, HISHIDA, SUZUKI, AND RUAN(GEN) "0, I 0.2 0.4 0.6 Time (ps) FIG. 4. Time dependences of the recombination luminescence (x) and of the self-trapped exciton luminescence (●) for liquid krypton. They are plotted so that the total counts for the two curves are equal to the ratio $L_1/L_2$. The long-time range (a) and the short-time range (b) are indicated. The solid lines represent fitted curves.

Figure 3.11 shows the scintillation pulse of electron, alpha and fission fragment interactions at zero electric field. It has also been demonstrated that the decay times are independent of the LET [22]. The scintillation of electron recoils in liquid xenon has a 45 ns fall time tail component at zero field that disappears when an electric field is applied. Figure 3.12 shows the electron recoil light pulse at zero and 4.5 kV/cm electric fields. This field dependence suggests that the tail is due to charge recombination. As this component is not observed

Table 3.3: Scintillation constants and intensity ratios of various particles in liquid xenon [71, 22].

<table>
<thead>
<tr>
<th>Particle</th>
<th>$\tau_1$</th>
<th>$\tau_3$</th>
<th>$(I_1/I_3)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Electron (1 MeV, 4 kV/cm)</td>
<td>2.2±0.3 ns</td>
<td>27±1 ns</td>
<td>0.05</td>
</tr>
<tr>
<td>Alpha ($^{210}$Po and $^{252}$Cf)</td>
<td>4.3±0.6 ns</td>
<td>22±1.5 ns</td>
<td>0.45±0.07</td>
</tr>
<tr>
<td>Fission Fragment ($^{252}$Cf)</td>
<td>4.3±0.5 ns</td>
<td>21±2 ns</td>
<td>1.6±0.2</td>
</tr>
</tbody>
</table>
for nuclear recoils (because of the much shorter recombination times), it serves as another mechanism for PSD in xenon. The recombination tail is also not observed for any type of recoil in liquid argon or neon, presumably for the same reason. The conversion of singlets to triplets (which is the lower energy excited state) by free electrons has been proposed as the mechanism behind the positive correlation between LET and $I_1/I_3$ [240, 22] – the electrons of lower LET events exist unbounded for a longer period and thus have a greater chance for converting states. The conversion of singlets to triplets is unlikely to occur because of the energy difference. This reaction has been observed in the gaseous state [241]. Section 5.3 describes the Case group’s study of pulse shape discrimination in a two-phase liquid xenon detector.

As previously mentioned, the Xe recoils deposit most of their energy into atomic motion which is not observed in liquid xenon detectors. The fraction of the total energy deposited in electronic interactions is typically known as the “nuclear quenching” factor. A theory formed by Lindhard [242] describes the appropriation of energy by incident ions into nuclear motion and electronic excitation. The relative scintillation yield at zero electric field between nuclear and electron recoils (usually with respect to 122 keV electron recoils), $L_{\text{eff}}$, is less than that of nuclear quenching as calculated by Lindhard because of two effects: “electronic quenching” and electron loss. As the scintillation decay times do not appear to be dependent on LET, the electronic quenching is presumed to occur before self-trapping, perhaps involving the loss of Xe*.

Hitachi proposes a mechanism for electronic quenching [21]:

$$\text{Xe}^* + \text{Xe}^* \rightarrow \text{Xe} + \text{Xe} + e^-.$$ (3.9)

One potential photon is lost in this process. A photon is also lost if an electron drifts away instead of recombining. The nuclear quenching (Lindhard) multiplied by electronic quenching factors gives the relative scintillation yield. The relative scintillation is very important as this value sets the recoil energy axis with respect to a gamma calibration line and consequently determines the expected WIMP event rate for a particular cross section.
The nuclear recoil energy based on $S_1$ is

$$E_r = E_e / \mathcal{L}_{\text{eff}} \cdot S_e / S_r,$$

(3.10)

where $E_e$ is the electron recoil energy scaled linearly from gamma peak and $S_e$ and $S_r$ are the electric field dependent scintillation yields relative to that at zero field for electron and nuclear recoils. Many groups have measured $\mathcal{L}_{\text{eff}}$ at energies relevant to WIMP searches. These experiments generally measure this value by irradiating a single-phase detector with a mono-energetic beam of neutrons. The recoil energy is constrained by tagging the scattered neutrons at different angles with a secondary liquid scintillation detector. Figure 3.13 shows the $\mathcal{L}_{\text{eff}}$ of these measurements including a theoretical prediction. Also included in the plot are the values predicted by Hitachi [21] and that determined by comparing XENON10 data with simulations [30].

### 3.5 Charge and Light Transport Properties

The electrons and photons that are generated by particle interactions must sometimes travel great distances before they are collected. The manner in which electrons and photons interact with the xenon is important as this affects the quality of the signal. Electrons diffuse as they drift through xenon and can be collected by charge impurities. Photons can also be absorbed by impurities. Due to its scintillation mechanism, the xenon is invisible to its own scintillation and thus the photons can travel far as long as the xenon is clean. In this section, we give a brief survey on the transport properties of electrons and photons in the xenon.

#### 3.5.1 Electron and Ion Mobilities in Xenon

As the electrons drift, they diffuse a distance $\sqrt{4Dt}$ where $D$ is the diffusion constant and $t$ is time. Figure 3.17 shows the energy dependence of the diffusion constants of electrons in liquid xenon and liquid argon [35, 243].
FIG. 3.13: The scintillation efficiency of nuclear recoils relative to 122 keV electron recoils.

The symbols correspond to this work: (□) Aprile et al. [21]; (△) Akimov et al. [29]; (×) Bernabei et al. [30]; (○) Ame´ndoa et al. [31]. The solid line represents the theoretical prediction by Hitachi [21]. This plot is from [24].

FIG. 8: Measured values $\mathcal{L}_{\text{eff}}$, the scintillation yield of nuclear recoils relative that of 122 keV gamma rays, for liquid xenon. Symbols correspond to this work: (□) Aprile et al. [21]; (△) Akimov et al. [29]; (×) Bernabei et al. [30]; (○) Ame´ndoa et al. [31]. The solid line represents the theoretical prediction by Hitachi [21]. This plot is from [24].

The work presented here represents a new measurement of $\mathcal{L}_{\text{eff}}$, which is also known by the proprietary names BC501A and NE213.

Acknowledgments

We express our gratitude to Dr. Steve Marino of the Columbia RaRAF facility for the beam time and his support throughout the measurements. We also wish to acknowledge the support from the University of Florida High-Performance Computing Center for providing computational resources.

As mentioned in the introduction, the uncertainty in $\mathcal{L}_{\text{eff}}$ at low recoil energies presents the largest systematic contribution to the WIMP exclusion limit. WIMPs of mass less than 8.8 $\times$ 10^3 GeV/cm^2 have been excluded at 99% confidence level.

The high-mass regime is relatively unchanged. In the literature [28] suggests a rise in the function going from 10 keV and below, this efficiency is lower than at 95 keV.

The work was supported by the Swiss National Foundation grants 20-118119 and 44-04-00596, and by the Swiss National Foundation grant 44-05825. This work was funded by National Science Foundation grant No. 04-04603, which is also known by the proprietary names BC501A and NE213.

We express our gratitude to Dr. Steve Marino of the Columbia RaRAF facility for the beam time and his support throughout the measurements. We also wish to acknowledge the support from the University of Florida High-Performance Computing Center for providing computational resources.
In the presence of an electric field, electrons and ions accelerate and achieve stable average drift velocities. The velocity of electrons in liquid and solid xenon increases with field until $\sim 1$ kV/cm as seen in Figure 3.14 [31]. For gas, the electron mobility is constant until about $E/N = 2 \times 10^{-2}$ Td ($1$ Td = $10^{-17}$ V cm$^2$) as seen in Figure 3.15. Note that the drift velocity is higher in liquid for the same $E/N$. Xenon ions and molecular ions ($\text{Xe}_2^-$) move much slower than electrons in liquid xenon as they are much heavier. Figure 3.16 is a plot of the mobilities of $\text{Xe}^-$ and $\text{Xe}_2^-$ in xenon gas as a function of the electric field.

### 3.5.2 Electron Capture

Impurities in the xenon can capture electrons. This decreases the size of the signal and introduces a position dependence. The latter effect is readily seen in two phase detectors, where the depth of the event determines how much liquid xenon the electrons must pass through before being extracted into the gas. There are three types of attachment processes [222]:

1) Radiative attachment

$$e + AB \rightarrow AB^- + h\nu$$  \hspace{1cm} (3.11)
The expansion contains the partial waves phase shifts $\delta_0(k)$, which are in a liquid, which are used for calculation of electron transport coefficients in dilute gas. Cross sections of electron scattering in liquid strongly differ from ones in a gas. Therefore the transport coefficients calculated for liquid differ from ones in dilute gas. The results of the calculations are presented in Figs. 4-6.

The effective polarization electron-atom potential has been used for calculation of electron transport coefficients in dilute gas. Cross sections of electron scattering in liquid strongly differ from ones in a gas. Therefore the transport coefficients calculated for liquid differ from ones in dilute gas. The results of the calculations are presented in Figs. 4-6.

\[ \delta_0(k) \]

Fig. 3.15. The electron velocity in liquid and gas xenon as a function of $E/N$ which is the electric field divided by the density (cm$^{-3}$) [32]. The points at the upper right corner are for xenon gas $E/N=10^{-17}$ V cm$^{-2}$. The momentum transfer cross section in the liquid is less than that in gas. The drift velocity is inverse proportional to the cross section and it is larger in the liquid for the same values of $E/N$. The calculation gives some decreasing of the drift velocity in the liquid for high fields. This is not verified by experiment where the drift velocity is saturated as a function of $E/N$ parameter for high fields. There is no Ramsauer minimum in the cross section of electron scattering in the liquid. Therefore the electron mobility in the liquid is a monotonic function of $E/N$ instead a maximum in the mobility in the gas. Fig. 5.

Heating of electrons by the electric field in the liquid begins for less values of $E/N$ parameter than in the gas. The mobility in the gas remains constant up to $E/N=0.02$ Td, while the mobility in the liquid begins to decrease at $E/N=10^{-4}$ Td. This is result of small cross section of electron scattering in the liquid compared with the cross section in the gas. This conclusion is verified by calculation of the mean energy of electrons in the liquid, Fig. 6.

The characteristic energy $\varepsilon_T$ is a quantity related to electron diffusion coefficient $D_T$ transverse to electric field $\varepsilon_T = eD_T/\mu$, where $\mu$ is electron mobility. This quantity has a dimension of energy and sometimes it is used for rough approximation in calculations. For example, the characteristic energy $\varepsilon_T$ in gaseous xenon is $9$ eV. In the liquid xenon the characteristic energy $\varepsilon_T$ is larger as a result of small cross section of electron scattering in the liquid. Therefore the electron mobility in the liquid is a monotonic function of $E/N$ instead a maximum in the mobility in the gas. Fig. 5.

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In argon lies 50 percent supports the assumption that the molecular argon purity ion.

Argon and xenon are shown to note that in these more difficult to form, \(^{18}\) is included in measuring the data are indicated by agreement with our higher \(E/\rho\) of Varney\(^{20}\) of the molecular ions, with Varney's measurements and experimental errors.

Reference \(26\).

Bernstein (private communication).


Fig. 10. Mobilities of \(\text{Xe}^+\) and \(\text{Xe}_2^+\) in xenon. The coding of the lines is the same as in Fig. 9.

21 I. B. Bernstein (private communication).

respectively, to the collecting electrode, and \(l_e\) and \(\tau_e\) are the electron drift length and drift time fit parameters, respectively, with \(l_e = v_e \tau_e\), where \(v_e\) is the electron drift velocity.

The electron drift length is obtained by fitting the \(S2\) of a gamma line to Equation 3.15. The rate of attachment (defined as \(k = (\tau_e n_i)^{-1}\) where \(n_i\) is the density of impurities) of electrons to \(\text{SF}_6\), \(\text{N}_2\text{O}\), and \(\text{O}_2\) as a function of electric field has been measured in \(\text{LAr}\) and \(\text{LXe}\) \(^{36}\) and is shown in Figure 3.18. The attachment rate to \(\text{SF}_6\) is much higher than that to \(\text{N}_2\text{O}\) and \(\text{O}_2\) at the electric fields measured – at 1 kV/cm, the difference is greater than two orders of magnitude. Between \(10^2\) to \(10^4\) V/cm, the attachment rate increases with field for \(\text{N}_2\text{O}\) but decreases for the other two solutes.

One of the early arguments against the development of commercial liquid xenon detectors was the lack of long-term purity. Now, purification can be achieved by passing xenon through various devices. One such product is the Oxisorb\textsuperscript{TM} purifier \(^{244}\), which contains a cartridge made of glass or aluminum that absorbs water and molecular oxygen. According to manufacturer specifications, water and oxygen can be cleaned to <30 ppb and <5 ppb, respectively, on input streams of 10 ppm. Xenon can also be purified with metal getters (e.g. Saes MonoTorr\textsuperscript{TM}). Spark purifiers have been applied to both liquid and gas xenon.
Figure 3.17: Electron diffusion coefficients in liquid xenon and argon versus electric field. The top axis is the electric field for liquid xenon case where $\rho = 3.0$ g/cm$^3$ [34]. The open circles are measurements by Derenzo [35].
Figure 3.18: The electron attachment rates to SF$_6$, N$_2$O, and O$_2$ in LXe as a function of electric field [36]
[245, 246], achieving drift lengths greater than a meter. This system operates on the same principle as that of the hot metal getters but with greater efficiency. The purifier works by creating sparks between two titanium electrodes, which create titanium dust with chemically clean surfaces that absorb impurities. In fact, it has been shown that the titanium dust can continue cleaning the xenon for years [247]. Purity is further enhanced by the UV light of the sparks, which breaks down complex molecules. Usually, the xenon is first purified with an Oxisorb cartridge and getter before being cleaned with the spark purifier. Spark purifiers have the advantage of being able to clean liquid xenon, allowing for a much higher purification rate. It has been shown that the level of purity achieved with chemical methods (Oxisorb™, molecular sieve, or metal getter) require 1–2 weeks of purification to reach the same level of purity achieved in 1–2 days with a spark purifier [245].

3.5.3 Transport of Charges Across Gas-Liquid Interface

Liberated electrons can be pulled out of the liquid and into the gas if they acquire enough momentum. The potential energy at the interface can be described by

\[
\begin{align*}
V_1(Z) &= V_0 - eE_1Z + A_1, Z < 0, \\
V_2(Z) &= -eE_2Z + A_2, Z > 0,
\end{align*}
\]

\[
A_{1,2} = \frac{-e^2(\epsilon_1 - \epsilon_2)}{4\epsilon_{1,2}(Z + \beta Z/|Z|)(\epsilon_1 + \epsilon_2)},
\]

where \(V_0\) is the ground state energy of the electron in the liquid; \(\epsilon_1\) and \(\epsilon_2\) are the dielectric constants of the liquid and gas, respectively; \(\beta\) is the thickness of the liquid-gas transition layer; and \(Z < 0\) is the liquid phase. Electrons with momentum greater than \(\sqrt{2m_e|V_0|}\) can leave the liquid. Figure 3.19 shows the potential energy versus \(Z\) for \(V_0 < 0\) and \(V_0 > 0\). Figure 3.20 shows measurements of the electron extraction efficiencies of argon and xenon in liquid and solid states [38].
(Table 1), \( \varepsilon_1 \) and \( \varepsilon_2 \) are the dielectric constants of the condensed non-polar dielectric and rarefied gas phase, the value of a cutting parameter \( \beta \) is about the thickness of the liquid–vapour transition layer which is probably a few times the condensed-phase thickness of the liquid–vapour transition layer—vapour transition layer interface is described by \( t_c = 1/(\alpha + 1/t_c) \). Then, \( \alpha = K = 1/[1 + t_c E_0] \) describes the process of heating the medium. These models are useful to describe the temperature dependence of the potential barrier in isooctane and liquid argon. In these media, the height of the potential barrier is comparable with the thermal energy of excess electrons which are in thermodynamic equilibrium with the medium. Electron emission \( (p_z > p_0) \) of high momentum electrons \( p_0 \) is possible only if these electrons have sufficient energy \( p_z \) for emission. The relaxation time \( \tau \) for emission of these high momentum \( p_z \) electrons is determined by the relaxation time for the thermal energy \( E_0 \) of excess electrons in the medium. Electrons in the high-momentum tail of the distribution function will have su\'cient energy for emission. The time-scale for emission of a large number of high-momentum electrons is given by \( \tau \) in the range of 10\( \sim 3 \) s.

In condensed media with so high \( |V_0| \) levels of purified activated electron emission from these media has not been observed. On the other hand, in these media the height of the potential barrier is comparable with the thermal energy of excess electrons in the medium. Then, these electrons may be emitted as thermal electrons. Otherwise, in dielectrics with high the thermal energy of excess electrons in the medium. Electrons in the high-momentum tail of the distribution function will have su\'cient energy for emission. The time-scale for emission of a large number of high-momentum electrons is given by \( \tau \) in the range of 10\( \sim 3 \) s.

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3.5.4 Proportional Scintillation

If electrons acquire sufficient drift velocity, they can excite the xenon gas and release scintillation. The number of proportional light photons created by electrons drifting in noble gases follows the empirical equation

\[ dN_{ph} = a \left( \frac{E}{P} - b \right) P \, dx, \]

(3.19)

where \( E \) [cm/kV] is the electric field; \( P \) [bar] is the pressure; \( dx \) [cm] is the distance traveled; and \( a \) and \( b \) are the fitted constants which are 70 photons/keV and 1 kV/cm/bar, respectively, for xenon gas [37].

3.5.5 Scintillation Propagation

Light attenuation can be described by an attenuation length \( \lambda_{att} \), which consists of two components, the absorption length, \( \lambda_{abs} \), and the elastic scattering length, \( \lambda_{sca} \). The relation
between the interaction lengths is

\[
\frac{1}{\lambda_{\text{att}}} = \frac{1}{\lambda_{\text{abs}}} + \frac{1}{\lambda_{\text{sca}}}.
\] (3.20)

In scintillation detectors, the \(\lambda_{\text{abs}}\) is much more important than the \(\lambda_{\text{sca}}\) as elastically scattered light can still be collected. Due to the unique scintillation mechanism of xenon, it is essentially transparent to its own scintillation. Impurities are responsible for light absorption. Two common light absorbing impurities in liquid xenon are water and oxygen. Figure 3.21 shows the absorption spectrum of 1 ppm of water and oxygen overlapped with the emission spectrum of xenon [248, 249]. The absorption length in liquid xenon has been measured to be greater than 100 cm in a 100 liter prototype developed for the MEG experiment [39].

Elastic scattering is dominated by Rayleigh scattering. The Rayleigh scattering length can also affect the final signal size – decreasing this length increases the total travel distance and thus raises the probability of absorption. The inverse of the Rayleigh scattering length is

\[
\frac{1}{\lambda_{R}} = \omega^4 \frac{\omega^4}{6\pi c^4} \left[ kT \rho^2 \left( \frac{\partial \epsilon}{\partial \rho} \right)_T^2 + \frac{kT^2}{\rho c_v} \left( \frac{\partial \epsilon}{\partial T} \right)_\rho^2 \right]
\] (3.21)
where \( \omega \) is the angular frequency of the scintillation, \( c \) is the speed of light, \( k \) is Boltzmann’s constant, \( T \) is the temperature, \( \rho \) is the liquid density, \( \kappa_T \) is the isothermal compressibility, \( c_v \) is the heat capacity at constant volume and \( \epsilon \) is the dielectric constant [250]. The scattering length has been calculated to be 30 cm [251] and measured to be 29 cm [252]. Equation 3.21 also shows how the spatial density fluctuations by temperature gradients can increase the amount of Rayleigh scattering. Note, that Equation 3.21 shows that the \( \lambda_R \propto \lambda_{\text{Xe}}^4 \), where \( \lambda_{\text{Xe}} \) is the scintillation wavelength. This means that increasing the wavelength can greatly extend the Rayleigh scattering length. This can be accomplished by coating the walls of the detector volume with a wavelength shifter or by doping the liquid.

As previously stated, most photons generated in the liquid tend to stay there due to the large mismatch in refractive indices of the liquid and gas. The refractive index of xenon gas is \( \sim 1 \). The refractive index of liquid xenon depends on wavelength (Figure 3.22) and density (Figure 3.23). The critical angle of total internal reflection at 170°C is \( \arcsin(1.0/1.69) = 36° \).
Figure 3.23: The index of refraction of liquid xenon as a function of temperature [42]. The discrepancy between this measurement and that by Barkov et al. [40], shown in Figure 3.22 may be due to a lower attenuation length in the latter measure.
Chapter 4

Case Detector

All the research and development described in the next chapter was performed with a cryostat that was constructed and tested at Princeton University in 2003 by Tom Shutt and the author. The cryostat and initial set of detectors were designed by Shutt. Many of the subsequent detectors and supporting structures were designed and constructed by the author.

This section gives a description of the Case cryostat, detectors and other supporting structures.

4.1 Case Cryostat

To obtain robust and consistent light and charge signals from the detector, the cryostat must provide a stable isothermal environment for the liquid xenon. In addition, the cryostat must also be reliable, easy to use and affordable. There are two common sources of cooling power. Cooling with liquid nitrogen is cheap and reliable as there are no moving parts. The detector can be cooled by extending a long rod or “cold-finger” between the liquid nitrogen and detector. The other option is to use a pulse-tube refrigerator (PTR), which is convenient but has a greater chance for mechanical failures (leaks, pump failures, etc.) thus necessitating a separate emergency cooling system. In a cold-finger cryostat, liquid xenon will simply freeze, during an electrical blackout as the heating units shut off. Another
Figure 4.1: Cold-finger cryostat designed by Tom Shutt.
disadvantage of the PTR system is the high upfront equipment cost.

We opted for a vacuum-insulated liquid nitrogen cooled cold-finger cryostat constructed primarily out of commercial parts. Figure 4.1 is a side view of the Case cryostat designed by Tom Shutt. The cryostat has a 22″ long, 2″ diameter aluminum rod as the cold-finger and three concentric cans: the outer chamber can, the radiation shield can and the Xe can. The use of modified commercial parts helps to lower the cost and allows for easier modifications. The outer can is constructed out of stainless steel 10″ ISO-LF parts and provides the vacuum insulation for the detector and aluminum rod. There are three 2-3/4″ Conflat ports with thin-walled tubing (to minimize heat loads) that extend from the Xe can. These three ports are each welded to the outer can with a bellows to eliminate stresses during thermal contractions.

The aluminum radiation shield can is bolted onto the aluminum base plate. The aluminum stick is attached to the bottom of this can, and bolted onto the other end of the stick is a copper tip that comes into direct contact with the liquid nitrogen. To conserve liquid nitrogen when running at higher temperatures, a thermal impedance (i.e. plastic disc) is inserted in between the bottom of the radiation shield and the aluminum stick. A 1/4″ copper tube loop is soldered onto the aluminum can with inputs and outputs at the bottom of the aluminum cold-finger. This loop serves as a conduit for liquid nitrogen, providing the >100 W cooling power needed for the initial cool down.

Attached to the copper piece at the end of the cold-finger is a pouch containing activated carbon (also called activated charcoal). The carbon improves the vacuum by adsorbing residual gas atoms through van der Waals attraction. Activated carbon has an exceptionally high surface area for adsorption – one gram has a surface area of $\sim 500 \text{ m}^2$.

The Xe can is a 316 stainless steel 8″ diameter, 6.6″ long Conflat half-nipple that holds the detector (as seen in Figure 4.2). The tube end is welded shut with a stainless steel plate. The center of this plate is partially bored out to a thickness of 1 mm to provide an entrance “window” for low energy gammas. A 1/4″ stainless steel tube is welded to the bottom of the can and serves as the outlet for the LXe. As the walls of the Xe can are fairly thin
Figure 4.2: Xe can, upside-down. The 1/4" stainless steel tube is the liquid xenon outlet tube. The apparent hole at the center of the bottom surface is the entrance “window” for low energy gammas. The copper block next to the tube is an enclosure for the platinum thermistor. Surrounding the walls of the can are many layers of copper foil which increase the temperature uniformity of the chamber.
and made of stainless steel which has low thermal conductivity, the opposite ends of the can are fairly decoupled and lead to non-uniformities in the LXe temperature. To increase the conductivity, several layers of copper foil ($k \sim 400 \text{ W/m/K}$) has been tightly wrapped around the walls of the can.

To introduce low energy gammas into the detector region through the window on the Xe can, an aluminum tube passes through openings in the outer vacuum can and radiation shield (as seen in Figure 4.1). This tube allows for the changing of radioactive sources without opening the cryostat.

Three 1\textquoteleft\textquoteleft diameter Vespel rods bolted onto the aluminum base plate and outer vacuum chamber support the weight of the internal structures. Vespel is used instead of metal because of its low thermal conductivity yet high strength. A stable cryostat temperature is maintained by a proportional-integral-derivative (PID) controller that powers a 50 W heater and monitors several platinum thermistors on the Xe can, radiation shield and cold-finger. Thin copper quad-twist cryogenic wires connect the thermometers to the controller. The standard four-wire measurement technique is used to eliminate the distortions from changes in the electric lines. The resistance of the Pt thermocouple is measured by passing a known current through two of the wires and measuring the voltage drop across the thermometer with the other two. There is no voltage drop across any of the voltage sensing wires as there is no current, thus the wire resistance is irrelevant. We observe a LXe temperature stable to within 0.1 K. Both the inner (Xe can) and outer vacuum spaces are monitored with ion gauges while pumping. To cover higher pressures, the inner and outer spaces are monitored with a Convectron (thermocouple) and Baratron (capacitance) gauges, respectively. Figure 4.3(a) and Figure 4.3(b) are plots of the temperature at the bottom of the Xe can and xenon gas pressure, respectively, over 20 days. The temperature rms was 16 mK over the 20 days and 5 mK over the first 5 days.

Care was taken to minimize the various sources of heat leakage. One source of heat is conduction through material which follows

$$P = \frac{kA}{L} \Delta T,$$  \hspace{1cm} (4.1)
Figure 4.3: The temperature at the bottom of the Xe can (top) and the pressure (bottom) versus time.
Table 4.1: Thermal conductivities common substances are room temperature [73].

<table>
<thead>
<tr>
<th>Material</th>
<th>Conductivity (W/m K)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Air</td>
<td>0.024</td>
</tr>
<tr>
<td>Aluminum</td>
<td>156</td>
</tr>
<tr>
<td>Copper</td>
<td>395</td>
</tr>
<tr>
<td>Diamond</td>
<td>2000</td>
</tr>
<tr>
<td>Gold</td>
<td>298</td>
</tr>
<tr>
<td>Polyimide</td>
<td>0.33</td>
</tr>
<tr>
<td>Silver</td>
<td>419</td>
</tr>
<tr>
<td>Stainless Steel</td>
<td>15</td>
</tr>
<tr>
<td>Water</td>
<td>0.59</td>
</tr>
</tbody>
</table>

where $k$ is the thermal conductivity, $A$ is the cross section area, $L$ is the length and $\Delta T$ is the temperature difference across the material. Table 4.1 shows the thermal conductivities of various materials at room temperature. The three 1.5″ stainless steel tubes, three Vespel rods and electrical lines introduce heat loads of about 2 W, 0.56 W, and several tens of mW respectively. The heat load by residual gas is negligible.

Another source of heat is black-body radiation. The total power of black-body radiation on the inner surface follows Stefan-Boltzmann law:

$$P = A\varepsilon\sigma(T^4 - T_0^4),$$

(4.2)

where $A$ is the inner surface area, $\varepsilon$ is the emissivity of that surface, $\sigma = 5.67 \times 10^{-8}$ J s$^{-1}$ m$^{-2}$ K$^{-4}$ is the Stephan-Boltzmann constant, $T$ is the temperature of the inner surface and $T_0$ is the temperature of the surrounding surface. Table 4.2 shows the emissivity of several materials. The emissivity depends on the surface temperature and finish. To reduce the radiative load, ten layers of aluminized mylar (super insulation) cover the sides and bottom of the radiation shield can, the Vespel rods, stainless steel bellows and the top of the aluminum base plate (as seen in Figure 4.4). The aluminized mylar has an emissivity of $\sim 5\%$. The radiative load without the mylar insulation is about 4.5 W. The ten layers cuts the radiation load by a factor of $\sim 10$. This insulation also lowers the heat load by lowering the mean free path of the residual gas particles.
Figure 4.4: The radiation can covered with 10 layer aluminized mylar. The Vespel rods and stainless steel bellows have also been covered with the insulation.
Table 4.2: Table of emissivities of common materials [74].

<table>
<thead>
<tr>
<th>Material</th>
<th>Temperature</th>
<th>Emissivity</th>
</tr>
</thead>
<tbody>
<tr>
<td>Aluminum, unoxidized</td>
<td>100°C</td>
<td>0.03</td>
</tr>
<tr>
<td>Aluminum, oxidized</td>
<td>199°C</td>
<td>0.11</td>
</tr>
<tr>
<td>Aluminum, heavily oxidized</td>
<td>93°C</td>
<td>0.20</td>
</tr>
<tr>
<td>Carbon, unoxidized</td>
<td>100°C</td>
<td>0.81</td>
</tr>
<tr>
<td>Copper, highly polished</td>
<td>38°C</td>
<td>0.02</td>
</tr>
<tr>
<td>Copper, black, oxidized</td>
<td>38°C</td>
<td>0.78</td>
</tr>
<tr>
<td>Glass, convex D</td>
<td>100°C</td>
<td>0.80</td>
</tr>
<tr>
<td>Glass, nonex</td>
<td>100°C</td>
<td>0.82</td>
</tr>
<tr>
<td>Iron, oxidized</td>
<td>100°C</td>
<td>0.74</td>
</tr>
<tr>
<td>Iron, unoxidized</td>
<td>100°C</td>
<td>0.05</td>
</tr>
<tr>
<td>Silver, polished</td>
<td>200°C</td>
<td>0.02</td>
</tr>
<tr>
<td>Stainless steel 301, polished</td>
<td>24°C</td>
<td>0.27</td>
</tr>
<tr>
<td>Stainless steel 301, polished</td>
<td>232°C</td>
<td>0.57</td>
</tr>
<tr>
<td>Stainless steel 303, polished</td>
<td>316°C</td>
<td>0.74</td>
</tr>
<tr>
<td>Stainless steel 316, polished</td>
<td>24°C</td>
<td>0.28</td>
</tr>
<tr>
<td>Stainless steel 316, polished</td>
<td>232°C</td>
<td>0.57</td>
</tr>
</tbody>
</table>

4.2 Gas System

The gas panel directs the flow of gas between the two storage bottles and between the bottles and the detector while purifying the xenon with a metal getter. Figure 4.5 is a photo of the bottles, panel and cryostat. Figure 4.6 is a schematic of the gas system. All components are connected with 1/4″ stainless steel tubing with Swagelok VCR fittings with metal gasket compression seals. These high performance seals have a maximum leak rate of $4 \times 10^{-9}$ std cm$^3$ s$^{-1}$. A regulator with a constant downstream pressure sits between the gas bottles and the panel. There are several Swagelok bellows valves for defining the flow path. The pressure at several points on the panel are measured with dial pressure gauges. The flow of gas is controlled by a calibrated mass flow controller. Also on the gas panel is a Monotorr getter purifier and gas pump. The panel was constructed inside a cleanroom to limit the entrance of dust. The 1/4″ tubing was cleaned in an ultrasonic bath with alcohol before installation. Before use, the gas panel and gas bottles were evacuated while baking at 100°C and 200°C, respectively. Before initial use, the entire system was checked with a helium leak detector.
Figure 4.5: Xenon bottles, gas panel and cryostat.
Figure 4.6: Schematic of the gas system.
The gas bottles are stainless steel Swagelok double-ended cylinders with one end welded shut. Welded onto the opposite end of each bottle is a $\sim$130 bar pressure rupture disc. The mass of the bottles are tracked with Ortec strain gauges.

Another way of cleaning the xenon and detector is by repeatedly recuperating all the xenon back into the bottles and then refilling the can. The xenon is extracted by cooling the bottles in liquid nitrogen. The regulator provides a constant 2.5 atm downstream pressure when filling from the storage bottles. Purification by recirculation is accomplished by pulling liquid xenon out of the can up the tube where it evaporates and pushing the xenon gas through the getter and back into the xenon can where it recondenses. The diaphragm gas pump drives the recirculation at a maximum rate of 3.3 liters per minute.

4.3 Detectors

Our primary goal in designing our detectors was to achieve high, uniform and stable light and charge collection. This requires highly reflective parts, a high quantum efficiency PMT placed in the detector volume and overall cleanliness. Details on the specific detectors are described in the next chapter along with their results. In this section, we give an overview of the various components.

4.3.1 Reflectors

Increasing light collection improves the signal resolution and threshold, both of which are key factors in the dark matter sensitivity. Highly reflective materials are placed around fiducial region to increase the photon collection. The primary reflector material used is PTFE (Teflon) which has been shown to have a reflectance to 175 nm photons ranging from 55% [43] to 95% [253]. The results observed in XENON10 suggest $>98\%$ PTFE reflectance. This large variation can be explained by two known sources of degradation: exposure to high intensity UV and the presence of light absorbing contaminants [254]. Figure 4.7 shows the reflectance of common reflectors below 220 nm [43]. In Section 5.4, we show the results of our attempt to improve light collection by use of a wavelength shifter and a reflector.
A simplified photon propagation Monte Carlo developed primarily by Tom Shutt was used to project the light collection efficiency and positional variation. Figures 4.8(a) and 4.8(b) are maps of the light collection efficiency of the bottom and top PMTs, respectively, of the Xed1b detector (which is described in greater detail in Section 5.2.1).

4.3.2 Wire Grids

Most of the wire grids used at Case and many of those used for earlier measurements at Columbia were designed and constructed by the author. We opted for wire grids over other forms of electrodes such as electroformed meshes (as used for XENON10) because they could be created in the lab thus affording a higher level of agility and lower cost. No field shaping rings were needed as the active volume geometry was always flat with a height of ∼1 cm and a diameter of ∼4 cm. Figure 4.10 is the electric field map of a system with two planes and two wire grids. These maps were created by Maxwell 2D to determine the

Figure 4.7: Reflectance of some common reflectors below 220 nm [43].

material called Spectralon.
Figure 4.8: Maps of the light collection efficiency of the bottom and top PMTs for light emitting from the active volume ($0.2 \leq z \leq 1.0$ cm) and region between the liquid level and anode grid ($1.2 \leq z \leq 1.4$ cm) where the proportional light is generated.
uniformity of the electric fields. The electric field of the active volume of our detectors were generally uniform to 5%. This high uniformity near the walls is due to the PTFE dielectric constant ($\sim 2$) closely matching that of LXe ($\sim 1.9$) and to the electrodes radially extending beyond the active volume.

**Selection of Wire Grid Materials**

All fine wires used for our experiments were purchased from California Fine Wire. In our prototypes, beryllium-copper (BeCu), gold-plated aluminum and tungsten wires were used. The reflectivities are likely to vary greatly with sample as features such as oxide layers and roughness can have a sizable impact (as they do for the emissivity). Because of these uncertainties, it is best to simply minimize the size and number of wires in the liquid. The wires in the liquid had a diameter of 40 $\mu$m or less. In the gas, the wire size is less important as most of the primary scintillation is retained in the liquid. To prevent electron multiplication in the gas (except in detector for gas gain measurement) these wires generally had a diameter of 125 $\mu$m. A uniform pitch (distance between adjacent wires) of 2 mm was
Figure 4.10: Electric field map of a two-phase setup with two wire grids (not visible) and two plates representing the PMT photocathode. The liquid level is between the two grids. The field is very uniform up to the edges because the grids extend into the PTFE walls and because the dielectric constants of the PTFE and LXe match well.

used for all grids. This dimension allowed for a large direct (>90%) transmission of light and charge.

Common ring materials are metals such as stainless steel and copper. Since we were interested in developing an individual wire charge readout system, we opted for 1/8” thick copper-plated Cirlex\textsuperscript{TM}, which is a thick polyimide adhesiveless laminate of the same material as Kapton. Cirlex is very strong and chemically resistant. Importantly, it has been shown by the Case group to not affect the electron drift length. The radioactivity of Kapton has been measured by the XENON collaboration to be 30-fold smaller than that of glass. Circuit patterns can be etched with laser or a milling machine. The former method is preferred as it leaves behind a much smoother service that is unlikely to trap dirt, but it is also more expensive and cannot be done in house. Figure 4.11 is pictures of a laser-etched Cirlex ring that was not used. The rings used in the measurements described in this thesis were etched with a milling machine.
Regular lead-tin solder was initially used, but was later replaced with silver-tin solder (5% silver) because of its higher melting point temperature. Although, the melting point of lead-tin solder is about 165°C, it softens at ∼120°C. Silver-tin solder melts at 220°C and has been tested to hold wires well up to at least 160°C. The higher operating temperature is important since it permits high baking temperatures. Furthermore, silver-tin solder has greater sheer strength and resistance to creep. Although not an issue in our prototype detector, silver-tin solder presumably has a lower radioactivity than lead-tin solder as it has much less $^{210}$Pb [255].

To ensure that the wires do not break or loosen when baking or cooling, the ring and wire materials are chosen so that the linear expansion coefficients ($\alpha$) closely match. The stress on a wire upon cooling or heating to temperature $T_C$ is

$$\sigma = E(\alpha_{\text{wire}} - \alpha_{\text{ring}})(T_0 - T_C),$$  \hspace{1cm} (4.3)$$

where $E$ is the Young’s Modulus; $\alpha_{\text{wire}}$ and $\alpha_{\text{ring}}$ are the linear expansion coefficients of the
Figure 4.12: Cirlex ring on the wire transfer frame. The 125 µm BeCu wires were soldered under a tension of 40 g to the 2 mm spaced copper pads of the circuit boards. At this stage, the next step would be to solder the wires onto the inner copper ring. Completed wire grids can be seen in Figure 4.13.

wire and ring materials, respectively; $T_0$ is the room temperature; and $T_C$ is the operating temperature. If the stress on a wire exceeds its yield point, then the wire may become permanently deformed. If $\alpha_{\text{ring}} > \alpha_{\text{wire}}$, then some tension needs to be applied when soldering the wire to the ring to ensure that the wire does not loosened when cooled. The necessary amount of tension is $\sigma A$, where $A$ is the cross sectional area of the wire. Note that Equation 4.3 is an approximation as it assumes that $\alpha$ is constant. For metals, $\alpha$ generally decreases as the temperature falls.
Figure 4.13: The Cirlex rings with 125 µm (left) and 40 µm (right) wires. The four on the periphery are the plates of the liquid level monitoring capacitor. Rings of this form were use for most of the measurements described in this thesis. In the center of each ring is a PTFE reflector.

Construction of Wire Grids

The wire grids are created by soldering tensioned fine wires onto copper coated Cirlex rings. The wires are first soldered onto the wire transfer frame (Figure 4.12). The frame has two G-10 boards with rows of copper pads spaced 2 mm apart. To ensure accurate pitch, each copper pad has a slight indentation across the center of each pad where the wire sits. The wires are soldered one at a time, under tension, to the transfer frame. A wire is tensioned across the frame by taping one end of the wire directly to the edge of one of the G-10 boards and taping the other end to a calibrated mass which hangs off the edge of the other G-10 board. After attaching the wires, the Cirlex ring is placed underneath the row of wires and the copper pads of the ring are covered with a generous amount of flux. The flux is added for two reasons. It removes the oxidation on the copper surface to ensure good bonding. It also acts as a surfactant to increase the fluidity of the solder which allows for easier spreading and increased smoothness of the final surface which is important for minimizing discharges. The wires are carefully soldered onto the inner copper ring and are then cut
with a razor above the outer copper ring. The sharp ends of the wires are covered with solder. The flux, which is fairly corrosive, is immediately removed by flushing with alcohol or acetone.

Before employing the wire grids, they undergo a temperature cycle test. In this procedure, the grids are heated on a hot plate at 150°C for 10 minutes, allowed to cool back to room temperature and then are placed in a liquid nitrogen bath. The wires are checked for tautness in the bath. This cycle is then repeated two more times and the grid is considered acceptable as long as none of the wires have loosened or snapped.

### 4.3.3 Photomultiplier Tubes

We used R6041 and R9288 Hamamatsu PMTs (as seen in Figure 5.29), which have the same compact profile with a 2″ window but with different quantum efficiencies (QE). The QE of the PMTs vary with wavelength and are 7% and 24% respectively, at 175 nm. The 12-stage dynode PMTs have a first dynode electron collection efficiency of 70%. They were designed to have low background and be able to operate in liquid xenon, which is crucial for light collection. The PMTs were positively biased to a voltage of 700–1000 V.

### 4.3.4 Liquid Level Measurement System

The liquid level can be measured by exploiting the difference in dielectric constants of xenon gas \(\epsilon_g \approx 1\) and liquid \(\epsilon_l \approx 1.87\). The capacitance changes as the xenon gas between any two electrodes are displaced by liquid. Two configurations were considered – a short tube with a thin central wire and parallel plates. We opted for parallel plates for several reasons. Firstly, capillary action in a narrow tube may cause inaccuracies in the capacitance measurement. Secondly, with the use of Cirlex rings, the parallel plates could be etched onto the rings and thus eliminate any ambiguity to the vertical position of the plates relative to the grids. In this section, we first derive the relation between the parallel plate capacitance and the liquid level and then describe the liquid level measurement system used in the Case detectors.
Figure 4.14: The liquid level (fraction filled) between the anode and gate wire grids versus volume while filling (above) and recuperating (below) at a constant rate. The erratic behavior at $\sim 0.7$ are likely due the liquid adhering to the high voltage leads attached to the anode grid. The large waves seen during recuperation are likely due to the boiling in the outlet tube.
Consider two parallel plates of surface area $A$ and gap distance $d$. The total capacitance is

$$\frac{1}{C_{\text{tot}}} = \frac{1}{C_1} + \frac{1}{C_2}$$  \hspace{1cm} (4.4)

as this system is essentially two parallel plate capacitors in series – one capacitor including the liquid ($C_1$) and a second one including the gas ($C_2$). These two capacitances are

$$C_1 = \epsilon_0 \epsilon_l \frac{A}{z},$$  \hspace{1cm} (4.5)  

$$C_2 = \epsilon_0 \epsilon_g \frac{A}{d - z},$$  \hspace{1cm} (4.6)

where $z$ is the height of the liquid level above the bottom plate, and $\epsilon_g$ and $\epsilon_l$ are the dielectric constants of the gas and liquid, respectively. The total capacitance of completely empty and completely filled capacitors are

$$C_{\text{empty}} = \epsilon_0 \epsilon_g \frac{A}{d},$$  \hspace{1cm} (4.7)  

$$C_{\text{filled}} = \epsilon_0 \epsilon_l \frac{A}{d}.$$  \hspace{1cm} (4.8)

Substituting these equations into 4.4, we obtain

$$C_{\text{tot}} = \frac{C_{\text{empty}}}{1 + \left(\frac{\epsilon_g}{\epsilon_l} - 1\right) \frac{z}{d}}.$$  \hspace{1cm} (4.9)

The dielectric constants can be replaced with the measured empty ($C_{\text{empty}}$) and filled ($C_{\text{filled}}$) capacitance values:

$$C_{\text{tot}} = \frac{C_{\text{empty}}}{1 + \left(\frac{C_{\text{empty}}}{C_{\text{filled}}} - 1\right) \frac{z}{d}}.$$  \hspace{1cm} (4.10)

We can take this equation and write the liquid level height, $z$, as a function of the total measured capacitance:

$$z(C_{\text{tot}}) = d \frac{C_{\text{empty}}}{C_{\text{filled}}} \frac{1}{C_{\text{tot}} - 1}.$$  \hspace{1cm} (4.11)

The $\sim 4$ mm gap between the gate and anode grids is monitored by three parallel plate capacitors built onto the Cirlex rings (as seen in Figure 4.13). There is one reference capacitor that uses the Cirlex as the dielectric. The pads have a width of $\sim 1$ cm and a
Figure 4.15: Liquid level (fraction filled) between the anode and gate wire grids at equilibrium. The gap between the anode and gate is 4 mm. The standard deviation of the fraction filled (liquid level) is 0.012 or (48 microns). The most likely explanation for the seesaw pattern is the gradually accumulation and sudden release of bubbles.

length of ∼3 cm. The signals are carried by Kapton coated copper wires. A Smartec UTI board measures the capacitances by a three signal technique that removes the dependence on signal gain and offset and nullifies the effects of instabilities in the processor clock cycle [256]. The system is sensitive enough to measure the minute level changes due to the pull of the liquid towards the grids when the electric fields are turned on. At the beginning of every run, the empty and filled capacitance values are measured. Then the detector is leveled by adjusting the three height adjustment bolts on the aluminum frame of the cryostat. Figure 4.14 shows the liquid level during filling and recuperation at a constant flow rate. The nonlinearity at ∼0.7 fraction filled is likely due to the liquid xenon interacting with the high voltage copper wires. The fluctuations are higher during recuperation because of the boiling occurring in the outlet line.

Immediately after filling the Xe can, there are typically large erratic swings in the liquid level due to falling droplets and escape of trapped bubbles. Several subsequent small fills and recuperations are required to achieve the desired level. After twelve or more hours, the liquid level always equilibrates to a see-saw wave with a period of several minutes (Figure 4.15). An explanation is the generation of bubbles by a hot spot at the joint between the bottom of the Xe can and the outlet tube. The bubbles slowly accumulate on a surface and
Figure 4.16: The Gaussian fit mean of $S_2$ versus liquid level (fraction filled) for 122 keV gammas at a stable liquid level setting. The spread in liquid level is due to the see-saw fluctuations (as seen in Figure 4.15). A Gaussian function was fit to the $S_2$ distribution of events in each liquid level bin. The error bars represent the statistical uncertainty to the Gaussian fits. There is an anti-correlation because an increase in liquid level decreases the electron path length in the gas. The amount of proportional scintillation follows Equation 3.19.

escape when they reach a certain size. This theory is supported by the disappearance of the waves upon heating the detector. Raising the temperature of the detector increases the gas pressure, which raises the energy barrier for nucleation. As previously noted, changes in the liquid level will alter the generation of proportional light. This see-saw fluctuation in the liquid level thus decreases the resolution of $S_2$. In a $^{57}$Co dataset with the average liquid level held constant, an anti-correlation between the liquid level and the $S_2$ of the 122 keV gammas can be seen. Figure 4.16 is a plot of the $S_2$ Gaussian fit means for several bins of liquid level. Recall that the amount of proportional light is dependent on the pressure,
electric field, number of electrons extracted and the distance the electrons travels in the gas (Section 3.5.4). Although the number of electrons extracted can increase with an increase in liquid level since the electric field is higher (if the gas field is above 10 kV/cm, then there is no change), the electron travel distance is the overriding factor. A liquid level correction to the proportional light was shown to slightly improve the resolution but since running the UTI board introduces noise, the liquid level is usually only measured before and after taking a dataset.

4.3.5 Electrical Lines

Commercial SHV hermetic connectors mounted in 2-3/4" Conflat flanges were used. These feedthroughs have a maximum voltage rating of 5 kV. Initially, single stranded bare 0.15" diameter copper wires were used to carry the high voltage from the feedthrough to the wire grids but due to the lack of insulation these were replaced with Gortex™ cable, which is
a multi-strand copper wire with plastic insulation, rated to several kV. There are several additional features to minimize electrical discharges. PTFE cylinders surround the point of connection between the feedthrough and Gortex cable (Figure 4.17). The connection between the Gortex wire and wire grid sits in the liquid where the break down electric field is very high. The PMT high voltage and signal lines are Kapton insulated coaxial cables designed for high vacuum systems and have a voltage rating of 1 kV. Several Bertan 375 high voltage power supplies were used to bias the grids. A system was later built to control these voltages with the DAQ computer. This allowed for the automated scanning across many electric field settings which significantly reduced the idle time of the detector. A Bertan 225 and an Ortec 556 are used to power the PMT’s. The Bertan 225 can be set to shut off at a particular current to prevent damage to the PMT. All the high voltage power supplies output voltages with 20–30 kHz noise which are removed with simple LRC filters.

### 4.4 Detector Cleanliness

To reduce the introduction of impurities, we used parts that were known or expected to be compatible with liquid xenon. New or well cleaned tools are used in machining soft materials (such as PTFE) to avoid embedding impurities. Before assembly, all parts of the detector are cleaned in an ultrasonic bath filled first with acetone and then with ethanol, both heated to 50°C. Parts that are particularly dirty such as freshly machined parts are first bathed in detergent and water. Next, the parts are flushed with de-ionized water and transferred to a clean room where they are left to dry. All parts are handled with clean disposable nitrile gloves. Care is taken to ensure very limited exposure of nitrile gloves to the acetone bath as nitrile dissolves in it. If possible, parts are also stored in a dry box before usage to remove some of the absorbed moisture. The detector is then assembled in a ultra-low particulate environment provided by a bench with a horizontal HEPA filtered blower (Figure 4.18). Before closing up, the detector is brushed with pressurized dry air or nitrogen to remove residue dust which are a source of electrical discharges.

The detector is usually baked at 50°C (limited by PMT) while being evacuated for
several days. After cooling down the detector, xenon is circulated for several days through
the detector (where it liquefies and absorbs impurities) and through a hot metal getter
purifier. For some experiments, the recirculation outlet was the tube welded to the bottom
of the Xe can. A PTFE outlet tube was later routed directly to the active volume to more
efficiently remove the impurities there. More information on this change can be found in
Section 5.4.1.

4.5 DAQ

Two types of signals are recorded – the fast nanosecond time scale light pulses from the
PMTs and the slower microsecond time scale charge signals from the spectroscopy amplifier.
Figure 4.19 is a simple diagram of data acquisition system (DAQ), which was programmed
in LabVIEW by Tom Shutt and Carl Eric Dahl. This section describes the PMT and charge

Figure 4.18: Clean bench with HEPA filter unit and blower.
4.5.1 Acquisition of PMT Signals

The PMTs signals are amplified near the signal feedthroughs with an SRS SR445A fast 5x amplifier to boost the signal-to-noise ratio. The amplified light signals are directed to a CAEN N401 Fan In/Out module. One set of signals is routed to the triggering system and another to the Acqiris cards. The triggering system consists of a Tektronics analog oscilloscope with a trigger out signal which is sent to the trigger input of the DAQ. For the two PMT setup, the DAQ is triggered by the signal from the bottom PMT as it collects more primarily light. The light signals are recorded by four 8-bit 500 MHz Acqiris digitizers. Two channels are needed for each PMT signal because of the great dynamic range in signal size and low resolution of the digitizers. The signals are stitched together in the data processing.

The primary light signal (S1) has a rise time of 5–10 ns and a fall time of about 50 ns. The pulse from the proportional light (S2) has a roughly Gaussian-shaped and a width of about 200 ns. The 500 MHz is well above the necessary sampling speed to capture the S1 and S2 pulses. Figure 4.20 is an example of the light signals from a 5 keV event in the bottom PMT of a Case two-PMT detector.

For some datasets, in particular the high statistics electron recoil datasets, we are primarily interested in low energy events. To conserve disk space, we sometimes implemented a high energy veto using gate/delay generators and multiple discriminators. For our setup,
a Philips 794 quad gate/delay module was used. Figure 4.21 is a schematic of this setup which functions in the following manner (see Figure 4.22): an S1 signal in PMT#1 (the PMT in the liquid) triggers the gate/delay generator to create a trigger signal, which is delayed by the maximum drift time, $\Delta t_{\text{max}}$. If there is an S2 from PMT#2 following the S1 that is within the maximum drift time and has a height that exceeds the upper threshold, a positive gate signal is created lasting for $\Delta t_{\text{max}}$. This is the veto signal that inhibits the third gate/delay from generating a trigger signal for the DAQ. Thus, if there is an S1 trigger and an S2 pulse with height greater than the threshold, the S1 trigger will be inhibited by the veto signal. This configuration also rejects events that have an S1 below the trigger threshold in conjunction to an oversize S2. In such a case, the S2 provides the primary trigger and trigger for the veto. To determine the veto trigger threshold, some calibrated
Figure 4.21: A schematic of the trigger setup with high energy veto. The signals are amplified by an SRS SR445A 5x amplifier. Copies are created with a CAEN N401 Fan In/Out. Three units on a Gate/Delay module are needed to create the veto logic. The system can be reverted to normal triggering (no veto) by disconnecting the line to the “inhibit” input.
Figure 4.22: This diagram briefly describes the operation of the high energy veto. The S1 creates a trigger signal that is delayed by the maximum drift time. An S2 that exceeds the threshold creates a inhibit gate signal lasting at least as long as the maximum drift time. This inhibit signal negates the delayed S1 trigger signal and the event is rejected.

Data is taken from which the S2 height threshold of the desired energy cut off is extrapolated. Note, that the S2 threshold must be generous to avoid cutting off the S1 pulse of same energy as there are large fluctuations between the sizes of the two pulses. Figure 4.23 is a histogram of S2 of two datasets, one without the veto and another with the veto at 30 keV.

4.5.2 Acquisition of Charge Signals

Direct charge measurements are made by integrating the charge with an Ortec 142AH pre-amplifier connected to the collecting electrode. Its feedback capacitance has been measured to an accuracy of <1% by multiple methods. Minimizing the input capacitance is crucial to lowering the noise so the shortest possible SHV cable (~4") is used between the HV feedthrough and the pre-amplifier input. The pre-amplifier signal has a rise time equal to the drift time of the electrons, which is usually about 5 μs. The integrated signal of the pre-amplifier is shaped and amplified by an Ortec 572 spectroscopy amplifier. The shaped signal is recorded by a National Instruments 6110 card, which has a sampling speed of 6
Figure 4.23: Histogram of S2 with and without the high energy veto. The high energy veto is set at \( \sim 30 \text{ keV} \).

MHz and voltage resolution of 14 bits. The DAQ is triggered by the shaped signal. The DAQ cannot be triggered by the pre-amplifier signal as it suffers from “flicker” noise.

### 4.6 Data Processing

After the data is written to hard disk, the files are compressed by creating new files that contain only the waveform of the pulses, omitting the sections of baseline that contain no useful information. Next, the basic quantities of the pulses are calculated and based on these values, the pulses are classified as S1, S2, charge pulses or noise. A framework was created
4.6.1 Data Compression

The light data requires a great deal of storage space. Two 10 µs PMT signals recorded at a speed of 500 MHz and resolution of 8 bits requires 80 kb or 10 kB of disk space. A million such events total 10 GB. Although the cost of hard disk space to store this data is not expensive, the transfer of such large files across the internet is problematic. For a full scale detector with scores of PMTs and a much longer maximum electron drift distance, the size of raw signal data is prohibitively large. Recognizing that most of the signal traces contain only small fluctuations around the baseline, these files can be compressed with fairly simple algorithms. Only the pulse waveforms, the pulse locations, and the baselines are saved to smaller files. With this information, a good approximation of the original waveforms can be
reconstructed. The pulses are identified by applying various filters. First, the baseline of the signal is subtracted and then the slow baseline shifts (∼1 μs) are removed with a high-pass filter. Next, the absolute value of the filtered signal is smoothed out with a Gaussian filter. The sections of the final filtered signal that are above a certain threshold are saved. Small segments preceding and following these sections are also saved. This procedure reduces the size of the data by a factor of ∼10. Figure 4.24 shows the original and filtered pulses.

4.6.2 Event Reconstruction

After the data is compressed, key characteristic values are calculated for the pulses and saved to file. These quantities include the pulse integral, height, rise time, fall time, start time, and the width at 10% and 50% of the maximum amplitude. Temporal moments such as the mean ($μ = [Σ_i t_i V(t_i)]/[Σ_i V(t_i)]$), skewness, and kurtosis are also calculated. Skewness or the third standardized moment, is $γ_1 = \frac{μ_3}{σ^3}$, where $μ_3 = [Σ_i (t_i - μ)^3 V(t_i)]/[Σ_i V(t_i)]$ and $σ = \sqrt{[Σ_i (t_i V(t_i) - μ)^2]/[Σ_i V(t_i)]}$. This quantity reflects the asymmetry of a distribution. Kurtosis is $γ_2 = \frac{μ_4}{σ^4} - 3$ where $μ_4 = [Σ_i (t_i - μ)^4 V(t_i)]/[Σ_i V(t_i)]$, and describes how peaked or flat a distribution is relative to the Gaussian distribution. A template pulse is also fit to the pulses to measure how closely the pulses conform to the expected pulse shape and to also provide a better estimate of the pulse height. This is important particularly for small charge signal pulses, where the height of the pulse is exaggerated by the noise fluctuations. The templates are average pulses.

In the next step, the pulses are classified based on the pulse shape. The most important quantities for discriminating light pulses are the pulse widths and rise times. For two-phase data, a good event is one that has a single S1 and single S2, within the maximum drift time expected. Since the Case prototype detectors had at most one PMT in the gas, there was no x–y position reconstruction capability. The z coordinate is provided by the drift time.
4.6.3 Computing Job Distribution

The compression of signals and calculation of pulse values demand a large amount of computational power so multiple computers are employed. With the data sets split among many files, there naturally exists a large number of discrete jobs that can be distributed among several computers. A multiple-slave system was setup in Matlab to handle job processing. The process starts with one of the computers creating a database file that categorizes which datasets that have been processed. Maintaining a database file is faster than requiring that each computer scour the hard drives every time that a job is needed. After the database file is created, all the computers take turns accessing the database file to check out jobs. Only one computer can access the database at any given time to prevent multiple computers from checking out the same job. The required data file is first copied over to the local hard drive. Once a computer finishes a job, it copies over the compressed data or pulse shape quantity file, deletes local copies of the files related to the job, and finishes up by marking the job as complete. Every hour or so, one of the computers checks to make sure that the database file is consistent with the actual processing state.
Chapter 5

Liquid Xenon Detector Research and Development

5.1 Introduction

In 1993, the ICARUS collaboration demonstrated discrimination between alpha and gamma ray events by measuring the primary and proportional signals of a two-phase liquid xenon time projection chamber [257]. An article published in 2003, showed the results of a two-phase xenon time projection chamber, which detected Xe recoils and distinguished them from electron recoils but lacked statistics and a clearly defined energy axis [258]. When the XENON collaboration formed in 2003, no two-phase liquid xenon detector had yet been employed for WIMP search. In fact, no one had determined the ultimate performance of the technology. Before starting on XENON10, several smaller prototypes were built by the Case Western group and our collaborators to gain experience with two-phase detectors, make calibration and discrimination measurements and explore alternative techniques for potential integration into the full scale detector. This chapter summarizes some of the basic research and development performed by the Case group, which is lead by Tom Shutt. The first two sections describes our measurements of discrimination by charge \( \log_{10}(S2/S1) \) and by shape of the primary scintillation pulse. The effectiveness of applying both princi-
ples simultaneously (as is done for two-phase liquid argon detectors) is also explored. The
following section describes the use of a wavelength shifter in an attempt to improve light
collection. The next section shows our operation of a detector with electron multiplication
running stably in two-phase. This demonstration is important because the use of a photo-
cathode (to convert primary scintillation photons into electrons) along with a direct charge
measurement system can be used in place of PMTs, which are expensive and radioactive.
The charge multiplication results are used to project the discrimination performance of a
detector with charge readout and PMT in the liquid. The final section shows the results of
a detailed measurement of the recombination fluctuations of 122 keV electron recoils, which
can be used to test recombination models.

5.2 Charge-Based Discrimination

When work began on the Xed prototypes in the summer of 2003, the ultimate discrimination
power of liquid xenon two-phase detectors had yet to be measured. Before observing Xe
recoils, we observed the recoils of $^{206}$Pb nuclei down to 20 keVr in our Xed1a detector,
which had only a single PMT in the gas. The $^{206}$Pb recoils were from alpha decays of
an internal $^{210}$Po source. Shortly afterward in November 2004, we measured the field and
energy dependence of the Xe nuclear recoil ionization yield and discrimination efficiency.
Some of these results are described in Section 3.4 and [19].

This section describes the results of a follow up measurement of the discrimination effi-
ciency with a higher light collection efficiency setup yielding a much lower energy threshold
and higher discrimination efficiency. Unforeseen structures in the two bands were observed
which fortuitously resulted in high background rejection down to 2 keVr. These results
established that this type of detector had the potential to set new cross section limits. The
next section describes the pulse shape discrimination using the same data.
Figure 5.1: Diagram of Xed1b detector which has two PMTs and five wire grids on three Cirlex boards. The 1/4" PTFE rod surrounded by the more attenuating LXe serves as the gamma collimator. Teflon reflectors line the inside of the active volume. The large PTFE block serves as filler material to lessen the amount of the LXe needed. Each PMT has a metal can surrounding the PMT at the base for electrical shielding.
5.2.1 Experimental Setup

The Xed1b detector had two Hamamatsu PMTs and five wire grid electrodes. Figure 5.1 is a CAD diagram of the Xed1b prototype used for the discrimination measurements. Figure 5.2 shows the detector from above with the PMT in the gas ("top") removed. Figure 5.3 is a view of the bottom of the detector with the lid of the stainless steel can shield removed. The base of the PMT in the liquid ("bottom") can be seen in this photo. The bottom PMT is a R9288 with a QE of 23%. The top PMT is a R6041 with a QE of 6%. The five wire grids are refered to, from top to bottom, as "top", "anode", "gate", "cathode", and "bottom". The top most Cirlex ring which holds the top and anode grids can be seen in Figure 5.3. Each wire grid had 20 wires spaced 2 mm apart. The top and anode grids consisted 125 µm gold plated aluminum wires. The gate, cathode and bottom grids were made of 40 µm BeCu. The cathode and gate, which defined the active volume, were separated by 1 cm. As with other detectors, a 1/4" diameter PTFE rod surrounded by LXe below the detector serves as a gamma collimator. The signal and high voltage pins on the PMT circuit board are electrically shielded by metal cans. One of these shields can be seen in Figure 5.3. Unlike Xed1a which had reflectors only in the liquid region, Xed1b had PTFE reflectors covering the active and gas regions. These reflectors not only help increase the light signal but also decrease the positional dependence. Figures 4.8(a) and 4.8(b) are light collection maps of the bottom and top PMTs, respectively, created by the photon propagation Monte Carlo. The detector is surrounded by a large PTFE block (see Figures 5.2 and 5.3) that extends to the walls of the Xe can and serves as volume displacement.

To measure the discrimination performance of a dark matter detector, we must find the discriminant response (charge, light, pulse shape, etc.) to background and WIMP events. With these measurements, we can then define the WIMP acceptance window which is the region in discriminant space that has a low rate of background but high rate of WIMP events. As previously mentioned, the dominant source of background are electron recoils from Compton scatters of high energy gammas (see Section 3.4 for information on gamma interactions with xenon). For background calibration, a hot gamma source is typically used.
Figure 5.2: View from above the detector with the top PMT removed. The white rings are the PTFE reflectors. The detector is surrounded by a volume displacing PTFE block. The pins are connections for the parallel plate capacitors.
Figure 5.3: This photo shows the bottom of the detector with the lid of the stainless steel PMT shield removed. Many components can be seen: the PMT shield can; the PMT base; the Kapton-coated coax wires that carry the high voltage and signal; the strip of Teflon that lines the inside of the PMT shield can and serves as an insulating barrier; the three 1/4" threaded PTFE support rods.
Figure 5.4: The $^{252}$Cf neutron energy spectrum [44]. The most probable and mean neutron energies are 0.7 MeV and 2.1 MeV, respectively.
rather than background in order to obtain a large number of events in a reasonable period of time. To calibrate the response to WIMPs, a neutron source is used. Neutrons with an energy of \( \sim 1 \) MeV mimic the response of \( \sim 10 \) keV WIMPs. High energy gammas were provided by a \(^{133}\)Ba source which was placed approximately 15” away from the center of the detector. A 25 \( \mu \text{Ci} \) \(^{252}\)Cf source provided neutrons with a falling exponential energy spectrum (as seen in Figure 5.4). The most probable and mean neutron energies are 0.7 MeV and 2.1 MeV, respectively [259]. Many more gammas are emitted than neutrons – the gamma emission rate of \(^{252}\)Cf is \( 1.8 \times 10^8 \, \text{s}^{-1} \mu\text{g}^{-1} \) versus \( 2.34 \times 10^6 \, \text{s}^{-1} \mu\text{g}^{-1} \) of neutrons – necessitating the use of a lead shield (see Figure 5.5). The \(^{252}\)Cf source was located 14” away from the center of the detector with 10” of lead in between. The lead placed behind the source reflected neutrons back towards the detector. The neutrons and gammas from these sources were not collimated.

At each drift field, a calibration (\(^{57}\)Co), gamma background (\(^{133}\)Ba) and neutron (\(^{252}\)Cf)
data sets were taken. Data was taken at five electric fields: 0.06, 0.52, 0.88, 1.95, and 3.96 kV/cm. A constant gas field of 10 kV/cm was set (5 kV/cm in liquid), providing full extraction of electrons into the gas [38]. The top and bottom PMTs were biased to +825 V and +875 V, respectively. The gamma background and neutron datasets contained $5 \times 10^5$ to $1 \times 10^6$ events.

The temperature of the liquid xenon was 187.4 K and the pressure was 1970 torr. The liquid level sat halfway between the anode and gate electrodes which was separated by 4 mm. The electron drift length was initially observed to be only several centimeters. The detector was purified by completely recuperating the xenon and refilling the xenon several times, passing the gas through the SAES MonoTorr getter in each direction of the cycle. Figure 5.6 shows the evolution of the electron lifetime as a function of recuperation/refill cycle. The detector eventually reached a drift length of $\sim 20$ cm which was constant throughout the rest of the run.

A general overview of the data processing is given in Section 4.6. A good event has one S1 and one S2 with a drift time within the range expected for the drift field. The pulse
integrals, $S1$ and $S2$, were corrected for depth dependence by using the response to 40 keV gammas from inelastic nuclear recoils in the $^{252}$Cf dataset and calibrated with the 122 keV gammas. The nuclear recoil energy follows Equation 3.10. For simplicity, we assumed a constant value of $L_{\text{eff}} = 0.2$ for the relative scintillation.

The measurement of the 122 keV gammas showed that the Xed1b detector had a primary light collection of 5.0 phe/keV at zero electric field for the bottom PMT. This corresponds to $\sim 50\%$ of the photons striking this PMT. The energy threshold was $\sim 2$ keVr which was set in the analysis. The charge threshold was $\sim 7$ electrons which was set by the trigger threshold. Only the PMT in the liquid is used for the $S1$ since the PMT in the gas collects $\sim 1/5$ as much light and has a far lower QE. The top and bottom PMTs collect roughly the same amount of proportional light. Either the top or bottom PMTs collect roughly the same amount of proportional light. For $S2$ of $\lesssim 300$ electrons, the bottom PMT is used because of its higher signal-to-noise ratio. However, for events of $\gtrsim 300$ electrons, the top PMT is used. Due to the high amount of light generated and the higher quantum efficiency and higher PMT bias of the bottom PMT, the $S2$ signal in the bottom PMT was highly nonlinear above several hundred electrons. For this reason, the charge signal of the 122 keV can only be measured by the top PMT. In order to calibrate the $S2$ of the bottom PMT in terms of charge, we examine events with $S2$'s that do not saturate the bottom PMT channel but are also above the threshold in the top PMT.

5.2.2 Results

As previously stated, the discrimination is rooted in the difference in linear energy transfer (LET), which manifests as $(S2/S1)_{\text{ER}} > (S2/S1)_{\text{NR}}$ (see Section 2.3.5). By convention, we use $\log_{10}(S2/S1)$ as the discrimination parameter. Figures 5.7(a) and 5.7(b) shows the $\log_{10}(S2/S1)$ versus energy of the $^{252}$Cf and $^{133}$Ba data sets, respectively, at 1.0 kV/cm. The nuclear recoil acceptance window is the region below the energy dependent mean of $\log_{10}(S2/S1)$. For all fields, the leakage is calculated for discrete energy bins with edges at 2, 6, 10, 20, 30, 40, 50, 60, 70, 80, and 90 keVr. The electron recoil leakage is defined as
Figure 5.7: Scatter plot of $\log_{10}(S2/S1)$ versus recoil energy of the $^{252}$Cf (top) and $^{133}$Ba (bottom) data sets at 1.0 kV/cm drift field. The background leakage events have been highlighted red. Also plotted are the band centroids of the electron (blue) and nuclear (green) recoils.
Figure 5.8: Histograms of and Gaussian fits to the $\Delta \log_{10}(S_2/S_1)$ of the electron and nuclear recoils peaks between 70 and 80 keVr at 0.88 kV/cm. $\Delta \log_{10}(S_2/S_1)$ is the $\log_{10}(S_2/S_1)$ subtracted by the mean of the nuclear recoils.

The discrimination efficiency is defined as one minus the leakage. We calculate two leakage quantities – actual and Gaussian. The actual leakage is determined by counting events inside and outside the acceptance window. For this quantity, a lower bound of $\mu - 2.5\sigma$ is included for the acceptance window to exclude events that are likely to be spurious. The Gaussian leakage assumes that the $\log_{10}(S_2/S_1)$ distribution is Gaussian and is calculated with two parameters: the Gaussian width of the electron recoil distribution and the distance between the means of the electron and nuclear recoil distributions. The electron and nuclear recoil energy dependent $\log_{10}(S_2/S_1)$ means are characterized by 10th order polynomials (these are represented by the lines in Figures 5.7(a) and 5.7(b)). These polynomials are calculated by iteratively fitting to the means of many overlapping energy bins, subtracting out the contour (“flattening”) and then repeating until curvature is omitted. Fitting Gaussian functions to the flattened bands provide a truer estimate of the band width since a shift
Figure 5.9: Actual and Gaussian electron recoil leakage versus energy at 0.88 kV/cm which is the drift field that provides the best performance. The error bars are statistical.

in \log_{10}(S2/S1) within an energy bin would otherwise lead to an overestimate. Figure 5.8 shows the histograms of the \log_{10}(S2/S1) of the $^{133}$Ba and $^{252}$Cf datasets with Gaussian fits between 70 and 80 keVr at 0.88 kV/cm.

Figure 5.9 is a plot of the actual and Gaussian leakages versus energy at 0.88 kV/cm, the drift field of the dataset that gives the best performance. In general, the actual leakage is much larger than the Gaussian leakage. We expected the actual leakage to be higher due to events that lose charge at the edges of the detector but can not be removed due to the lack of x-y position information. XENON10 and the ZEPLIN detectors confirmed this. Figure 5.11 shows the \log_{10}(S2/S1) fit means of both recoils as a function of energy. There are several key features in this plot. Both bands turn up at lower energies because there are fewer ions with which the electrons can recombine with [60]. In our prior measurement with Xed1a, which had a much higher threshold of 20 keVr, the sharp upturn was apparent in the nuclear recoil band but it was unknown whether this feature would also be present in the
electron recoil band. Had the electron band not shifted upward, the discrimination would have been far worse. This feature has been confirmed by XENON10 and other experiments [13, 207]. The $\log_{10}(S2/S1)$ of the electron recoils appear to be much more field dependent than that of the nuclear recoils which is a reflection of the difference in ionization density. A proposed model for the band turn-up at low energies will be part of a future publication by our group [260] and in another thesis [60].

In summary, we remeasured the discrimination efficiency of two-phase liquid xenon detectors to lower energies with a higher light collection efficiency prototype, finding a rejection efficiency that is $>99.5\%$ below 20 keVr. This observation matches the results of XENON10. The field and energy dependence of the band structures are also explored.
5.3 Pulse Shape Discrimination

5.3.1 Introduction

As previously mentioned, the measurement of the scintillation light alone can also provide
discrimination, albeit at a lower efficiency than charge-based discrimination (see Section
3.4). Dark matter detectors using pulse shape discrimination (PSD) with NaI \[183\], high
pressure xenon gas \[261\], single phase liquid xenon, argon and neon \[12, 195, 262, 263, 264\],
and two-phase argon \[203, 265\] have been employed or are currently under development, but
thus far there has been no exploration of PSD in a two-phase liquid xenon time projection
chamber. In a two-phase detector, \(\log_{10}(S2/S1)\) can be used to clearly separate the electron
and nuclear recoil events for better measurement of their respective pulse shapes. This
section describes the measurements of the electric field and energy dependence of PSD in
such a detector with the same data as that used to explore the charge-based discrimination
described in the previous section. Also explored is the use of the prompt/total pulse shape
quantity in conjunction with \(\log_{10}(S2/S1)\) to further reduce the background beyond that
achievable by either discriminant alone. Lastly, PSD performance is projected to larger
5.3.2 Prompt/Total Method

The PSD performance was examined for various pulse shape quantities such as the width at half-max, integral/height, prompt/total fraction and a quantity called “multi-bin,” which is a generalization of the prompt/total fraction to more than two bins (described in [200]). The prompt/total and multi-bin quantities provide equal discrimination efficiency in liquid xenon. Here, we describe only the results of prompt/total which is defined as

\[ f_p = \frac{\int_{t_i}^{t_f + t_{\text{window}}} V(t) \, dt}{\int_{t_i}^{t_f} V(t) \, dt}, \]

(5.1)

where \( V(t) \) is the signal voltage, \( t_{\text{window}} \) is the prompt window size, \( t_i = t_0 - 50 \) ns, and \( t_f = t_0 + 300 \) ns, where \( t_0 \) is the time at which the pulse reaches 10% of the maximum amplitude on the leading edge. The optimal \( t_{\text{window}} \), depends on the PMT, electronics, and detector size. For Xed1b, it is 10 ns and for the Monte Carlo pulses of larger detectors as described in the next section, \( t_{\text{window}} = 36 \) ns is used.

5.3.3 Pulse Shape Monte Carlo

A pulse shape Monte Carlo was developed to gain a better understanding of the \( f_p \) distribution and to project the PSD performance to larger detectors by incorporating a photon arrival time distribution. To simulate the S1 signal, we need the photon emission (excimer decay) probability distribution function (PDF), the photon arrival time PDF, PMT gain distribution, and photoelectron response waveform. The PMT gain distribution is the single photoelectron distribution from the photoelectron calibration dataset. The photon emission PDF is obtained by deconvolving the photoelectron response from the average S1 pulse. Deconvolving the arrival time is unnecessary due to the small size of the detector as confirmed by the photon propagation Monte Carlo (described in Section 4.3.1). The average pulses are calculated by lining up the pulses according to their \( t_0 \) and calculating the mean voltage for each time bin. This was done for electron and nuclear recoil events.
Figure 5.12: Average pulses (top) and corresponding photon emission probability distribution functions (bottom) for 78 keVr electron and nuclear recoils at 0.58 kV/cm as determined with the Case prototype.
Figure 5.13: The distributions of photon arrival time to the bottom PMT array for various scales of LUX (top) and for events at the top and bottom of LUXx4 (bottom). The top plot represents an average of the active volume.
Figure 5.14: Electron and nuclear recoil pulses in various detectors generated by the Monte Carlo. The pulses become less distinct with increase in detector size as the fluctuations in arrival time dominate.

in 2 photoelectron bins. The PMT response to a single photoelectron was obtained in the same manner from the photoelectron calibration dataset. Figures 5.12(a) and 5.12(b) show the average S1 pulses and corresponding emission PDF for 78 keVr electron and nuclear recoils at 0.06 kV/cm.

The photon arrival time PDF for larger detectors was obtained with the previously mentioned photon propagation Monte Carlo. In addition to simulating Xed1b (to confirm the assumption of instantaneous photon arrival times), a next generation two-phase liquid xenon detector called the Large Underground Xenon Detector (LUX), and a detector with four times the dimensions of LUX (LUXx4) were also simulated. Nominal LUX active volume dimensions of 49 cm diameter and 54 cm height were used. LUX and LUXx4 have active masses of roughly 300 kg and 20,000 kg, respectively. Five aluminum meshes of 0.02 normal opacity are placed in same manner as in Xed1b. The arrival time distributions were simulated for 50 different \((r, z)\) coordinates. Figure 5.13(a) shows the photon arrival time...
Figure 5.15: Plots of $f_p$ versus energy of the nuclear (red) and electron (black) recoil events (top) of the $^{252}$Cf data set and of the electron recoils events (bottom) of the $^{133}$Ba data set at 0.06 kV/cm. The leakage events of the $^{133}$Ba data set have been highlighted red. The nuclear and electron recoil centroids are indicated by the dashed green line and solid blue line, respectively.
distribution for several scales of the LUX detector. Figure 5.13(b) shows the arrival time distribution for photons emitted near the top and bottom of the active volume of LUXx4. This plot shows that photons emitted from the top of the active region generally strike the bottom PMT array much later and with a significantly wider distribution in arrival times than those emitted from the bottom. Figure 5.14 shows the average electron and nuclear recoil pulses for the different sized detectors as created by the pulse shape Monte Carlo. The electron and nuclear recoil pulses become noticeably less distinct as the detector size is increased. At each electric field and for each of the three detectors, 24,000 light pulses of each recoil type were simulated between 0 and 100 keVr.

5.3.4 Results

Figure 5.15(a) is a scatter plot of $f_p$ as a function of recoil energy for the $^{252}$Cf dataset. Figure 5.15(b) is a similar plot of the $^{133}$Ba data set with the leakage events highlighted.
Figure 5.17: Gaussian leakage of $f_p$ as a function of recoil energy at 0.06 kV/cm. Also included are the Monte Carlo predictions for larger detectors.

The nuclear recoil acceptance window is the region above the energy dependent nuclear recoil mean as indicated by the green dashed line. The leakage is calculated in a fashion that is identical to that used for $\log_{10}(S_2/S_1)$ in the previous section (Section 5.2.2). The 10th order polynomial fits to the band mean can be seen in Figures 5.15(a) and 5.15(b). Figure 5.16 shows the histograms of $f_p$ between 55 and 60 keVr for the $^{133}$Ba and $^{252}$Cf (excluding the ER events) at 0.06 kV/cm.

Figure 5.17 shows the electron recoil leakage as a function of energy at 0.06 kV/cm which is the drift field that provides the lowest leakage. Also included in this plot are the Monte Carlo predictions of the detector-averaged leakage for LUX and LUXx4. We see that the PSD is very weak for LUXx4 and thus would probably never be used in such a large detector except for regions of the detector near the PMTs. The discrimination efficiency of LUX is $\sim 90\%$ for recoil energy $>70$ keVr and steadily decreases with decreasing energy. The discrimination efficiency in Xed1b is $\geq 95\%$ at energies above 50 keVr but is much lower below 20 keVr, which is the energy region of interest for dark matter detection. The error
Figure 5.18: The energy dependence of the electron and nuclear recoil $f_p$ band separation (top) and fit sigmas (bottom) of the electron recoils at 0.06 kV/cm drift field. The error bars represent the statistical error.
bars of the leakage include only the statistical fit uncertainties. The existence of systematic shifts in the pulse shape between the electron and nuclear recoil datasets would manifest as a shift in the \( f_p \) of the electron recoil as they are present in both datasets. Several changes could lead to a change in pulse shape: the introduction of an impurity that captures and emits light with a decay time, a drastic shift in the geometry or a change in the electronics. None of these changes are likely to have occurred and in any case, the fit mean of the electron recoil pulse shape of the \(^{252}\text{Cf}\) and \(^{133}\text{Ba}\) datasets agree within error.

Figures 5.18(a) and 5.18(b) show the band separation and electron recoil band width as a function of energy, respectively. These figures reveal that the rise in leakage with decreasing energy is a result of the convergence of the two bands and an increase in the band width. The convergence of the \( f_p \) of the electron and nuclear recoils can be explained by the convergence in LET (see Figure 3.9). Also included in Figure 5.18(b) are the Monte Carlo data band widths, which match well with the data indicating that band width is dominated by statistical fluctuations. The slightly higher width in the actual data suggests that some other intrinsic component (e.g. recombination fluctuations) that is not included in the Monte Carlo is apparently making a small contribution to the band width. The data shown in these plots can be used to estimate the discrimination performance of a detector with higher light collection. We obtain the intrinsic component to the width as the difference between between the Monte Carlo and observed widths. The band width at energy \( E \) of a detector with a factor \( x \) higher light collection is the sum (quadrature) of the Monte Carlo width at \( x \times E \) plus the intrinsic width at \( E \). Here, we see that although an improvement in the light collection would result in a large reduction in the band width below 20 keVr, the leakage is unlikely to drop meaningfully as the band separation is near zero.

Figure 5.19 is a plot of the field dependence of the leakage between 60 and 80 keVr. The leakage increases as the electric field is increased. The band width appears to be independent of the electric field (Figure 5.20(b)) as expected since the Monte Carlo shows that the width is dominated by statistical fluctuations which do not change with electric
field. The trend in the leakage can be explained by the change in the $f_p$ mean (Figure 5.20(a)). The $f_p$ of nuclear recoils appears to be field independent while that of the electron recoils possesses a strong field dependence below $\sim 1$ kV/cm, similar to the behavior of the electron recoil charge yield. The change in the tail is attributed to the removal of electrons from the event site. As the electric field is increased, the electrons that would have otherwise recombined at longer time scales are swept away by the electric field. The presence of electrons may also be changing $f_p$ by converting singlets to triplets [240, 22]. On the other hand, the charge yield of the nuclear recoils changes little with field [19] as does the pulse shape. This mechanism suggests that a correlation should be observed between mean $f_p$ and $\log_{10}(S_2/S_1)$, as the latter parameter is a measure of charge yield. A positive correlation is observed for the mean values across electric fields (Figure 5.21), however, the correlation at an electric field is roughly zero. The correlation is either nonexistent or is masked by the much large statistical fluctuations.

The electron recoil leakage can be reduced by changing from the usual nuclear recoil window defined by $\log_{10}(S_2/S_1)$ to a two-dimensional window in $\log_{10}(S_2/S_1)$-$f_p$ space.

Figure 5.19: Leakage versus electric field between 60 and 80 keVr using $f_p$. The error bars represent the statistical uncertainty.
Figure 5.20: The Gaussian fit means for electron and nuclear recoils (top) and fit sigmas for electron recoils (bottom) between 60 and 80 keVr as functions of the electric field. The error bars represent the statistical uncertainty.
Figure 5.21: Plot of the mean $f_p$ versus the mean $\log_{10}(S_2/S_1)$ between 60 and 80 keVr at five different electric fields. The error bars represent the statistical uncertainty.

(“2D window”) as the discriminants are not redundant. For simplicity, the acceptance window is defined by a diagonal straight line cut through the center of the nuclear recoil cluster, retaining the 50% acceptance for nuclear recoils. For projections to larger scale detectors, the fits to the $\log_{10}(S_2/S_1)$ of the real data is used in conjunction with the fits to $f_p$ of the Monte Carlo data. Figures 5.22(a) and 5.22(b) are scatter plots of $f_p$ versus $\log_{10}(S_2/S_1)$ of the $^{133}$Ba and $^{252}$Cf data between 70 and 80 keVr at 0.06 kV/cm, with the $\log_{10}(S_2/S_1)$ and 2D windows, respectively. The circled dots are the electron recoil leakage events. Figure 5.23 shows the leakages with the two different windows in Xed1b and the projected performance for LUX. The results for LUXx4 are omitted because the improvement is minuscule. The error bars are statistical. In Xed1b, the reduction in Gaussian leakage below 30 keVr is 16% and for LUX, 5%. Note that these are detector averages. The fluctuations in photon arrival time of events near the bottom of the detector are smaller than average (Figure 5.13(b)) thus the performance of this technique in such regions would be better than that shown in Figure 5.23.
Figure 5.22: Scatter plots of $f_p$ versus $\log_{10}(S_2/S_1)$ between 60 and 80 keVr at 0.06 kV/cm with the $\log_{10}(S_2/S_1)$ (top) and 2D window (bottom) cuts shown. Both the $^{133}$Ba (black) and $^{252}$Cf (red) points are plotted with respect to the nuclear recoil mean. The blue circle highlights ER leakage events. The dashed lines are the $2\sigma$ ellipses.
In summary, we have measured the field and energy dependence of pulse shape discrimination in a two-phase liquid xenon detector. By using $\log_{10}(S2/S1)$, we were able to cleanly separate the electron and nuclear recoils for a more robust measurement of the respective pulse shapes. With the aid of a pulse shape Monte Carlo, we find that the width of the pulse shape bands are dominated by statistics; other processes like recombination fluctuations contribute little. We observe no apparent correlation between the fluctuations of pulse shape and $\log_{10}(S2/S1)$ at an electric field. A significant increase in light collection is unlikely to meaningfully increase the PSD efficiency at low energies relevant to dark matter searches. Although, the discrimination by pulse shape is poor, it can be used with the charge measurement to reduce the leakage beyond that achievable by either method alone. Finally, the pulse shape Monte Carlo is used to project the performances of PSD and 2D-window to larger detectors.
5.4 Use of a Wavelength Shifter in a Two-Phase Xenon Detector

The discrimination performance of liquid xenon dark matter detectors greatly depends on the light collection efficiency as the statistical fluctuation of the primary light signal is a limiting factor at recoil energies of several keVs [266, 267, 60]. XENON10 and the Case prototypes with two PMTs achieve a light collection efficiency of about one photoelectron (pe) per keVr (slightly field dependent). The use of wavelength shifters have been used in LXe scintillation detectors [266]. Shifting the emission spectrum to longer wavelengths can boost the light signal in several ways. First, the reflectance of PTFE and other materials are generally higher for photons of higher wavelength. As previously stated, the reflectance at 175 nm has been reported to be between 55% and 95% (see Section 4.3.1 on discussion on reflectors). The reflectance at 350 nm has been measured to be 98.5% [45]. Spectralon, a proprietary form of PTFE created by Labsphere, has a reflectance that rises from 98% at 300 nm to >99% above 400 nm [45]. The spectral reflectance of Spectralon, PTFE and other reflectors to photons of wavelength <220 nm can be seen in Figure 4.7. The reflectance to photons of wavelength >300 nm is shown in Figure 5.24. Secondly, the Rayleigh scattering length increases with wavelength as $\lambda^4$ (see Equation 3.21) [251]. Extending the Rayleigh scattering length decreases the total travel distance of the photons, which decreases the likelihood for absorption by impurities in LXe (assuming that the wavelength shifted photons are not more readily absorbed by impurities). Finally, the quantum efficiency of photomultiplier tubes typically increases with wavelength. For example, work at the Jefferson Lab concluded that the deposition of p-terphenyl (p-TP) onto the glass windows of Burle 8854 PMTs increased their quantum efficiency in the blue range by about 40% [268].

A group exploring the benefits of using wavelength shifters in dark matter detectors has successfully operated a NaI detector with p-TP wavelength shifter at 160 K [269]. ICARUS T600 [205], ArDM [205], and WARP [203] are three argon detectors that use the wavelength shifter tetraphenylbutadiene (TPB). TPB has also been used in a liquid neon detector [72].
Degradation of the diffuse reflectance of Spectralon under low-level irradiation

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Published 7 February 2003
Online at stacks.iop.org/Met/40/S212

Abstract
Spectralon™ is a well known diffuse reflectance standard. Although ageing effects of Spectralon under high-level VUV and air UV irradiation have been reported, Spectralon is assumed to be stable under the normal radiation conditions used for calibration in standard facilities. In this paper, ageing effects of Spectralon under low-level irradiation and in the dark are presented. It has been found that ageing is mainly affected by four parameters: wavelength of radiation, irradiance level, radiant exposure (dose) and storage time of the samples. Therefore, the use of Spectralon™ as a transfer standard in international comparisons and maintenance of reference values is discussed.

1. Introduction
Spectralon™ is a well established reference material for diffuse reflectance standards [1]. Besides good Lambertian scattering behaviour, the advantage over other comparable materials such as pressed BaSO₄ tablets, opal glasses or ceramic tiles is that the reflectance is very high even in the UV spectral region due to the base material polytetrafluoroethylene (PTFE) and can easily be used to prepare integrating spheres. The hardness of the surface is durable enough to be cleaned and to even use it in facilities with soft surface touch.

In earlier publications [4–7], a strong decrease of the reflectance of Spectralon samples from more than 95% to below 30% at 300 nm has been reported [4], especially in the short wavelength region due to high-level VUV or solar-like radiation. Nevertheless, Spectralon was assumed to be stable under the normal radiation conditions used for calibration of reflectance standards in the visible and UV spectral region. In addition, one publication [8] showed evidence for the ageing of Spectralon used in integrating sphere wall material.

This report presents changes of the spectral reflectance of Spectralon samples observed after irradiation by different sources, causing ageing effects even at low levels used for calibration. This degradation is due to the level of irradiation and wavelength of the radiation and depends on the production lot and the storage time.

Wavelength shifters are required in argon and neon detectors because of their low wavelength emission (128 nm and 77 nm, respectively). We are unaware of any published results on the use of a wavelength shifter in a two-phase xenon detector. In this section, we report the results of running such a detector with vacuum-deposited p-TP.

5.4.1 Detector

Figure 5.25 is a CAD diagram of the Spectralon detector, Xed1f. Figure 5.26 is a photo of the detector with the top reflector removed. This detector consisted of five wire grids and a single 2" diameter Hamamatsu R9288 PMT in the liquid. The high voltage and signal connections at the PMT were shielded with a stainless steel cup. The wire grid
Figure 5.25: CAD diagram of the Xed1f detector with Spectralon reflectors.
just below the liquid level (gate) consisted of 40 µm diameter BeCu wires; the other four had 125 micrometer diameter gold-plated aluminum wires. The active volume was defined by the cathode and gate grids and had a depth of 1.0 cm and a diameter of 3.83 cm. Five Optical-Grade Spectralon reflectors – five rings and one circular plate – surrounded the active volume. Since the reflectance of Spectralon is very sensitive to impurities such as oils, we had the manufacturer of Spectralon, LabSphere, machine the parts. We chose Spectralon over Teflon even though the literature shows identical reflectance at high wavelengths because the Spectralon is likely generated under more controlled conditions and thus is less likely to show the large variations in reflectance that has been observed for Teflon at 175 nm. Its reflectance is a function of thickness (as seen in Figure 5.27) – the reflectance is 0.935 at 1 mm and rises to ∼99% at 5 mm and levels off. All five reflector pieces have a thickness of at least 5 mm. Spectralon is extremely hydrophobic as it has a
2.4 Reflectance Properties of Thin Sections of Spectralon

The reflectance of Spectralon decreases with decreasing thickness over most of the spectrum. Thin sections of Spectralon, less than 4 mm, may be doped with barium sulfate to maintain high reflectance and diffuse properties. The figures below illustrate the reflectance properties of thin sections of Spectralon and barium-sulfate-doped Spectralon.

![Figure 10](image)

**Figure 10**: The reflectance of Spectralon and Spectralon doped with barium sulfate to 325 nm photons as a function of thickness [46].

![Figure 11](image)

**Figure 11**: Reflectance properties of thin sections of Spectralon and barium-sulfate-doped Spectralon at 555 nm.

![Figure 12](image)

**Figure 12**: Reflectance properties of thin sections of Spectralon and barium-sulfate-doped Spectralon at 850 nm.
p-TP has the best record in applications for noble gas detectors because of its reasonable stability in vacuum, low hygroscopy, chemical inertness, and extreme radiation hardness. Several studies have proven that p-TP does not contaminate pressurized xenon gas and helium with electronegative impurities and has been demonstrated to have a >90% quantum efficiency for converting 175 nm photons into 340 nm ones [272]. Spectralon has a reflectance of 98.5% at 340 nm [45]. Figure 5.28 shows the emission spectrum of p-TP [47].

5.4.3 Vacuum Deposition of Wavelength Shifter

There are several common ways of depositing wavelength shifters. One method is to spray coat the items with a mixture of p-TP dissolved in toluene. Another is by polymer matrix coating. In this technique, the wavelength shifter and a plastic is dissolved in toluene, and the mixture is coated onto the parts and allowed to dry. The most complicated method is by evaporation in a vacuum chamber. We opted for this technique because it affords the best purity, uniformity and control of thickness [273]. A standard vacuum thermo-deposition
system was used to deposit 0.5 mg cm$^{-2}$ of p-TP on the Spectralon reflectors and the window of the PMT. Before deposition, the p-TP was baked at 50$^\circ$C in vacuum to remove absorbed water and other impurities. The sublimation began at a pressure of 2 $\times$ 10$^{-7}$ torr and temperature of 160$^\circ$C. The coated pieces were allowed to cool down in the nitrogen atmosphere before removal. The thickness of deposited layers was estimated by measuring the change in the mass of the p-TP and calculating, by geometry, the fraction of p-TP that lands on the Spectralon pieces. Before usage, the coated pieces were stored in a dry box to prevent absorption of water which can lower the quantum efficiency of p-TP [274].

5.4.4 Light Collection Efficiency

The expected number of photoelectrons generated for an event of energy $E$ is

$$N_{pe} = \left(\frac{E}{w_{ph}}\right) \cdot LCE \cdot QE,$$

(5.2)

where $w_{ph}$ is the energy required to generate one photon (15.2$\pm$0.3 eV at zero field from the calibration measurement described in [275]), $LCE$ is the light collection efficiency, and $QE$ is the quantum efficiency of the PMT. The number of photons emitted is $E/w_{ph}$. The number of photons expected to strike the photocathode is $N_{pe}/QE$. With $N_{pe}$, we have the $LCE$ of the detector. The $N_{pe}$ is just the $S1$ of the 122 keV electron recoils divided by the mean single photoelectron response and thus only two datasets were needed. Both zero field datasets were taken with the detector at 175 K and liquid level above the surface of the top Spectralon disc reflector. Only events within the space 4–6 mm above the cathode are used. With this cut and the collimation by the Teflon rod, only events at the center of the active volume are considered. The peaks are fit with Gaussian functions. This setup had a photoelectron conversion efficiency of 6.5 pe/keV at zero drift electric field. This value has a statistical uncertainty of less than 1%.

To correctly assess the affect of using a wavelength shifter, a similar single PMT detector with no wavelength shifter should also have been tested but the experiment was terminated due to the very poor charge collection caused by the p-TP (see next section). The closest comparison that can be made is with the performance of the bottom PMT of the two PMT
Xed1b detector running in two-phase. The bottom PMT in the prototype with Teflon reflectors was 5.0 pe/keV. The light collection efficiencies of the Teflon and Spectralon setup were 47±1% and 57±1%, respectively – the Spectralon setup has 21% higher light collection than the Teflon setup. Note that part of the difference in the light collection efficiency is due to the difference in the detector configurations. In the two-PMT setup, the photons that leave the liquid are unlikely to reenter it and be collected by the bottom PMT.

There are several potential sources of systematic error in the light collection efficiency values. The first is the nonlinearity in the PMT signal response to larger S1’s due to the current output limitations of the PMT. This effect was determined to be negligible for the signals of this experiment. The second is the possible change in quantum efficiencies when cooled to cryogenic temperatures. Having observed no change in the QE’s by changing the temperature from 165 to 185 K, we assumed that these values are constant with temperature. The errors given for $LCE$ are due to uncertainties in $w_{ph}$. A homemade light propagation Monte Carlo written in MATLAB was used to determine the reflectance of the Teflon given the observed light collection and determine the expected light collection efficiency of the wavelength shifter plus Spectralon setup. Simulated photons reflect diffusely off non-metal surfaces and shift in wavelength before reflecting off surfaces covered with p-TP. Photons are propagated until absorption. Refractive indices of 1.68 [42] and 1.51 (given by the PMT manufacturer) were used for the LXe and PMT glass, respectively, and values of 30 and 100 cm were used for the Rayleigh and absorption scattering lengths [222, 39]. Reflectance values of 0.40 and 0.35 were used for scattering wavelength shifted photons off BeCu (assumed to be the same as Cu) and Au, respectively [276]. The reflectance of BeCu and Au below 200 nm could not be found and likely varied greatly depending on the surface quality. Since varying the reflectance of BeCu at 175 nm from 0 to 1 shifts the $LCE$ by less than 1%, a value of 0 was arbitrarily chosen. The reflectance of Au is 0.3 near 200 nm [276]. Assuming that the Au reflectance at 175 nm is 0.3±0.3, the Monte Carlo can match the $LCE$ of the two PMT Teflon setup with Teflon reflectance of 0.85±0.05. This
The value falls within the range of previous observations (see Section 4.3.1) but is much lower than that seen in XENON10, which saw a reflectance of >98%. The presence of embedded light-absorbing impurities or degradation by UV light may be reasons for the difference [254]. There is also the possibility that the Monte Carlo underestimated the reflectance by omission of certain detector details such as small gaps between reflectors rings which would have otherwise absorbed light and decreased the expected number of collected photons. A $LCE$ of 67±2% (error from the uncertainty in Au reflectance) was expected for the Spectralon setup based on the Monte Carlo. A post-run inspection revealed that some of the p-TP had migrated to other parts of the detector, which could account for the lower than expected $LCE$. Figure 5.29 shows the p-TP-coated PMT before and after the experiment, showing that large amounts of the p-TP had been removed. The omission of small gaps between reflector rings in the Monte Carlo may also have contributed to the discrepancy in light collection efficiency.

5.4.5 Electron Lifetime

To measure the effects of impurities on the charge collection in LXe, $^{57}$Co datasets were taken with the detector running in two-phase configuration at 175 K. The electric field in
Figure 5.30: Electron drift length (dots) and temperature (dashed line) versus time. There is clearly an anti-correlation between the drift length and temperature that can be explained by the temperature-dependent solubility of p-TP. The error bars are statistical errors of the fit.

the active volume was 1 kV/cm. The field underneath the liquid surface was 5 kV/cm which provided full electron extraction. The liquid level was set halfway between the anode and gate grids. The electron lifetime and drift length were extracted as described in Section 3.5.2. The electron drift length has routinely been measured to be greater than \( \sim 100 \) cm in our prototypes without p-TP. However, in this detector, the electron drift length stabilized at \( \sim 0.4 \) cm after more than a week of recirculation. Various attempts at purifying the xenon did not improve the situation. This included recirculating the xenon through the metal getter purifier at different flow rates and modifying the detector to improve the efficiency by which the liquid in active volume of the detector was passed through the purifier. This final modification involved adding a PTFE outlet tube directly to the active volume (Figure 5.31). Furthermore, no vacuum leaks were found in the detector or associated plumbing and a subsequent run with a detector devoid of the wavelength shifter showed excellent purity, thus the equipment was likely not at fault. The impurity problem was likely related to the
presence of the wavelength shifter.

Influence of Electric Field Strength

At a fixed recirculation flow rate and xenon temperature, we measured the electron lifetime as a function of the drift electric field. Electron life time drops with increasing electric field strength (as seen in Figure 5.32) corresponding to an electron attachment rate that increases with field like $N_2O$ (see Figure 3.18).

Influence of Temperature

A potentially more important insight into the impurity issue was revealed by the observation that the drift length did not further decrease over time when recirculation purification was turned off. This is in marked contrast to the common situation early in a run when outgassing appears to be the dominant source of impurities. Moreover, in static conditions with no recirculation, the drift length was inversely related to the temperature by an apparently fixed relationship that was reversible over several cycles of temperature (as shown in Figure 5.30). This strong and reversible correlation between the drift length and temperature suggests that the p-TP is dissolved in the LXe at a level determined by a temperature-dependent solubility. Molecules of p-TP may be acting as acceptors themselves or be facilitating their capture by other impurities by slowing down their drift velocity.

Compounds with higher molecular weight are likely to have lower solubility in LXe. From this point of view, it looks reasonable to repeat the experiment with heavier members of p-N-phenyl family such as p-quaterphenyl (p4P), p-quinquephenyl (p5P), and p-sexiphenyl (p6P); Table 5.1 shows their properties. In fact, p4P has been used as an effective extreme-ultra-violet wavelength shifter in the past (see, for example, [277]) and has demonstrated emission and absorption spectra similar to that of p-TP [278]. Also, p-N-phenyls have demonstrated high chemical stability and high luminescence quantum yield in the blue range [279]. All of them can be deposited by evaporation with the sublimation temperature elevated for heavier members. These organic substances attract attention because of their
Figure 5.31: The PTFE outlet tube connected to the active volume of the detector which is shown here upside down in the clean bench. The liquid xenon was originally pulled out from underneath the detector but in an attempt to better clean the impurities out of the active volume, the PTFE outlet tube was connected coupled to this region.
Figure 5.32: The electron lifetime versus the electric field.

Table 5.1: Properties of p-N-phenyls. p3P = para-terphenyl; p4P = para-terphenyl; p5P = p-Quinquephenyl; p6P = p-Sexiphenyl

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<tr>
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<th>p3P</th>
<th>p4P</th>
<th>p5P</th>
<th>p6P</th>
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<tbody>
<tr>
<td>Molecular Weight</td>
<td>230.1</td>
<td>306.41</td>
<td>382.50</td>
<td>458.59</td>
</tr>
<tr>
<td>Melting Point (°C)</td>
<td>212</td>
<td>&gt;300</td>
<td>381</td>
<td>475.16</td>
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</table>

potential use in organic light-emitting diodes, organic field-effect transistors, and solar cells [280]. As previously mentioned, some argon and neon detectors use TPB. However, these detectors operate at a far lower temperature and thus it is not clear that TPB would work in liquid xenon.

In summary, we have demonstrated the use of a wavelength shifter in a LXe detector. Our first choice of p-TP as the wavelength shifter results in unacceptable levels of charge loss, but other promising wavelength shifter materials with potentially lower solubility in LXe remain an attractive option.
5.5 Charge Multiplication

5.5.1 Introduction

Measuring proportional scintillation is an excellent way of achieving a charge measurement with low threshold and high resolution. However, this standard method requires many photomultiplier tubes which are a large source of radioactivity. The Hamamatsu PMTs used in our prototypes are also expensive (∼$2000 each), fragile, and require several months to manufacture. Charges can be directly measured with a charge sensitive amplifier attached to the anode electrode. However, directly measuring the charges of several tens of electrons is a very difficult task given the typical amount of electronics noise. The charges can be measured if the electrons are multiplied through ionization avalanches in high electric fields around thin wires or other micro structures. This method can lessen the number of PMTs or eliminate them completely. In the latter case, the primary light can be measured by having a cathode plate covered with a photocathode material (such as CsI) for converting the primary light into electrons. In a detector with photocathode, the proportional light of one set of electrons can produce a set of even more electrons. This mechanism can grow and eventually evaporate the photocathode. A solution is to include a fast high voltage...
switch to turn off the extraction field for a certain amount of time after the necessary pulses have been recorded. If the necessary quantum efficiency for the photocathode cannot be achieved, the bottom PMT array can be retained to measure the primary and proportional scintillation while two perpendicular sets of wire grids in the gas would provide position sensitivity and a second measurement of the ionization. The ArDM is a two-phase liquid argon dark matter detector under development that has such a configuration [205].

This section reviews prior research performed with noble elements, gives a basic overview of the physics of charge amplification, and presents the results of measurements performed in our two-phase xenon prototype.

### 5.5.2 Previous Work on Electron Multiplication

Electron multiplication was demonstrated 30 years ago in liquid xenon using thin wires with a thickness of several microns [281]. The fragility of the fine wires motivated the exploration of microstrips. One such experiment used microstrips of 8 µm wires on an insulating borosilicate substrate, which observed a gain of about 10 in LXe [282].

Electron multiplication with a triple Gas Electron Multiplier (GEM) structure has been demonstrated in two-phase detectors using argon, xenon, and xenon doped with methane (CH$_4$) by Bondar et al [49]. Figure 5.33 is a plot of the electric fields and electron tracks of an avalanche around a GEM hole. Methane is added as a quenching gas, which usually increases the maximum gain and stability by suppressing feedback (by blocking photons from initiating secondary avalanches). However, its use in a dark matter detector would be problematic as methane would also suppress the S1 signal. Figure 5.34 shows the gain in pure and CH$_4$-doped xenon versus the voltage difference across the GEM. The detector obtained a reproducible ∼5000 gain stable to 6 hours with saturated argon vapors at a pressure of 0.76 (85 K) and 1 atm (87 K). However, such high gains could not be achieved with saturated xenon vapors. The detector obtained a maximum gain of ∼200 with pure xenon at 0.88 atm (163 K) and 0.99 atm (165 K) and was able to operate stably for a half-hour with a gain of 80. The addition of 2% CH$_4$ to saturated xenon vapors did not
We studied the possibility to use the two-phase Xe detector with the CH4 gas, namely methane. In that case, the detector was prepared in the bottle at room temperature. After the maximum dissolution level of CH4, it was to be 1.9% [23], the CH4 concentration in the two-phase mode might be taken as approximately 2%. Figure 12 shows gain–voltage characteristics for a two-phase Xe + CH4 mixture in the two-phase mode, when there is no liquid condensate at the bottom. The latter is close to the gaseous mode, when there is no liquid condensate at the bottom. At cryogenic temperatures, it is shown that the maximum gain can easily exceed $10^4$, while in the two-phase mode the maximum gain does not exceed $10^3$, as in two-phase Xe. Thus, the performance of two-phase Xe avalanche detectors, based on GEM structures, can be operated in an electron-avalanche counting mode for the first time. Emission and stability characteristics obtained for these avalanche detectors are rather promising. In Ar, two phase avalanche detectors with pure xenon were achieved with gains, reaching 5000, were obtained in a GEM detector with pure xenon operating in two-phase Xe. The characteristics obtained in two-phase liquid-gas Ar and Xe detectors are rather promising. The former also could not operate for periods greater than a half-hour, presumably because of the condensation on the multiplication structures. A GEM detector with pure xenon was able to operate in two-phase with a gain of 25 for several hours [283]. It is not clear how this detector was able to run stably for such an extended period while the others could not.

Both the primary light and charge must be measured for background discrimination. Measuring the primary scintillation requires a photocathode such as CsI or NaI coated with CsI in pure xenon has been successfully demonstrated at room temperature. Figure 5.35 shows the charge signal including the pulses induced by the scintillation (equivalent to a triple-GEM of a three GEM structure [49]. The presence of CsI on the primary light signal is enhanced by the generation of proportional ionization. The primary light signal is also shown.

A gain of 80 under irradiation with $\beta$-particles. During this period the three-GEM operation was relatively stable and the gain variations observed (of a factor of 2) were correlated to the vapor pressure variations (of about 5%). It should be remarked that this result is rather correlated to the vapor pressure variations (of about 5%). The presence of CsI on the primary light signal is enhanced by the generation of proportional ionization. The primary light signal is also shown.

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the collecting anode, one gets a prompt signal corresponding to photoelectrons produced by the primary light emission, followed by a later pulse with a delay corresponding to the drift time of ionization electrons in the detection gap. The detector can be used to determine the time of conversion of neutral radiation, and therefore the penetration depth, potentially allowing correcting the parallax error intrinsic in thick gaseous devices. The principle was demonstrated long ago with the so-called scintillating drift chamber, making use of an external photomultiplier to detect the primary light flash [5]. Recently, the same approach is used in dual-phase detector, where the primary scintillation in liquid xenon is exploited in conjunction with the measurement of drift time of ionization electrons extracted from the liquid into the gas phase, where the charge can be multiplied and detected [6, 7]. Various types of gas devices with internal CsI photocathodes are also described in Ref. [8], also reporting the observation of primary scintillation for low energy X-ray sources, albeit with rather low efficiency.

To investigate this possibility, we have mounted a small size (3 cm² active) multi-GEM detector in a vacuum-grade vessel having signal and HV feed-through; the vessel could be evacuated before gas filling. All measurements have been done in pure xenon at 1 bar. X-rays are converted in a 5 mm thick drift gap, followed by four multipliers at 2 mm spacing; the top electrode of the first GEM is coated with CsI. Drift field values up to 3 kV/cm could be reached in the conversion gap, permitting one to largely enhance the photon yield by secondary emission. Signals are collected on the last electrode, a single pad with the size of the multipliers. For the present study, the detector has been exposed to 5.9 keV X-rays, emitted by an internal collimated 55Fe, and 22 keV X-rays from an external 109Cd source.

Fig. 2 shows examples of prompt scintillation signals, followed by the main charge pulse, for two values of drift field (200 and 2000 V/cm). In the first case, the early pulse is due to primary scintillation, while the ramping signal at higher field is due to the insurgence of secondary scintillation, produced by inelastic collisions of drifting electrons (existing data indicate a threshold for its onset at around 800 V/cm in xenon at 1 bar [9]). Notice also the small after-pulses; they are probably due to secondary electron extraction from the photocathode by positive ion produced in the first GEM [10], and probably set a limit to the maximum gain of the structure.

Dividing the amplified signals into two channels, sent to separate discriminators set at very low and high threshold, we recorded the time difference between the early and main pulse. Fig. 3 shows the resulting distributions for the 5.9 and 22 keV sources at low drift fields; the dashed line corresponds to time distribution measured for a CsI scintillator.
In this section, a brief overview of the physics of electron multiplication and a derivation of the first Townsend coefficient which is the number of ionizations per unit length over a distance \(dr\) is

\[
dN = N_0 \alpha dr,
\]

where \(\alpha\) is the first Townsend coefficient which is the number of ionizations per unit length of the mean free path between ionizations. The first Townsend coefficient is determined by excitation and ionization cross sections and on various transfer mechanisms. The first Townsend coefficient is

\[
\alpha = \frac{dN}{dt} = \frac{dN}{dx}
\]

where \(dN/dt\) is the number of ionizations per unit time and \(dN/dx\) is the number of ionizations per unit length.

In off-line data analysis, a simple algorithm was used to look for within a certain time interval prior to the S2 signal. Fig. 6 shows peak delay with respect to the S2 trigger at a gain of 2500 and amplifier shaping time of 0.5\(\mu s\). A more complete overview can be found in [286].

Due to its complexity, \(\alpha\) must be measured for each gas mixture. Figure 5.38 shows \(\alpha\) for several gases. The first Townsend coefficient for drift fields of 0.25 and 0.61\(kV/cm\). One can see that the time delay between S1 and S2 depends on the drift field and is larger for lower fields. At a drift field of 0.25\(kV/cm\) its most probable value is \(-2.4\mu s\), which corresponds to electrons drifting through the liquid and gas gaps.
particles in the two-phase Ar avalanche detector with CsI photocathode at a gain of 2500, drift field \( E(LAr) = 0.25 \text{kV/cm} \) and amplifier shaping time of 0.5 \( \mu \text{s} \).

Fig. 7 shows the distribution of events in the plane of \( S2 \) versus \( S1 \) amplitudes, at a gain of 2500. One can see that most events are of the \( S1 + S2 \) type where both \( S1 \) and \( S2 \) signals are observed. Note that they are correlated (proportional) to each other. On the other hand, some events have an \( S2 \) signal alone, presumably due to a reduced geometrical acceptance for scintillation detection, for example, in those cases where ionization is produced in the corner of the active area. Here the \( S1 \) signal is just a noise. A negligible amount of events are due to \( S2 \) treated as \( S1 \); in Fig. 6 these are contained in the tail of events with a positive time delay.

Fig. 8 shows amplitude spectra of the \( S1 \) and \( S2 \) signals. The \( S2 \) spectrum has a single peak corresponding to the high-energy component of the \( \beta \)-particle spectrum in Fig. 5. The low-energy component was disregarded since it was under the trigger threshold. The \( S1 \) spectrum has also a peak, of similar shape as the \( S2 \) peak. In addition it has a tail at lower amplitudes, corresponding to electronics noise. This tail reflects the events of the \( S2 \) alone type in Fig. 7.

In Figs. 7 and 8, (top scale), the \( S1 \) and \( S2 \) amplitudes are expressed in units of the initial charge prior to multiplication, i.e. in photoelectrons and electrons, respectively. The appropriate scale calibration was carried out using the gain value and the amplifier calibration at a given shaping time. It should be noted however, that the accuracy of this calibration procedure was not too high: the corresponding scale could be off by a factor of 1.5 either way. More details will be presented in a more elaborated paper [12].

As one can see from Fig. 8, the number of photoelectrons in the \( S1 \) peak is about \( N_{pe} = 30 \). This corresponds to 30 photoelectrons per 600 keV of deposited energy. Hence, the figure of merit (FOM) expressed in units of the initial charge prior to multiplication is given by

\[
\text{FOM} = \frac{N_{pe}}{E_{ph}} = \frac{30}{600} = 0.05.
\]

It should be noted that this estimation gives only a lower limit of the photon detection efficiency, in particular, due to the effects of photoelectron backscattering to the CsI photocathode itself, the photon collection efficiency of the CsI photocathode itself, the photon detection efficiency is substantially reduced, i.e. by a factor of 50, to the fact that the scintillation light yield in liquid Ar, \( \alpha \), has for the first time been demonstrated in the two-phase avalanche mode. The detection efficiency is estimated to be of the order of 10\(^{-3}\) per 600 keV of deposited energy.

3. Conclusions

The performance of a two-phase Ar avalanche detector with CsI photocathode was demonstrated for the first time in liquid Ar, operated in saturated vapour above the liquid phase and a xenon atom. Integrating Equation 5.3 from where the avalanche starts \( r_{min} \) to the surface of the wire \( r = r_{max} \) gives

\[
I_{\text{ph}}(E) = \frac{4}{3} \pi \int_{r_{min}}^{r_{max}} \rho(r) r^2 dr,
\]

where \( \rho(r) \) is the density of the active area at the surface of the wire (\( r = r_{max} \)). The performance of a two-phase Ar avalanche detector with CsI photocathode was demonstrated in the two-phase Ar, 3GEM detector [51]. The gain value and the amplifier calibration at a given shaping time were under the trigger threshold. The \( S1 \) spectrum has also a peak, of similar shape as the \( S2 \) peak. In addition it has a tail at lower amplitudes, corresponding to electronics noise. This tail reflects the events of the \( S2 \) alone type in Fig. 7.
where $E_{\text{min}}$ is the field at $r_{\text{min}}$. Substituting in Equation 5.5 for $E$ gives

\[
M = \exp\left(\int_{E_{\text{min}}}^{E(a)} \frac{\lambda \alpha(E)}{2\pi\epsilon_0 E^2} dE\right).
\]  
(5.8)

To proceed further with this integral, a functional form of the $\alpha$ is needed. In the region of $E/P$ between $10^2$ to $10^3$ V cm$^{-1}$ torr$^{-1}$, $\alpha$ is roughly linear with $E$ (as seen in Figure 5.38), and thus we can substitute $\alpha = \beta E$ into Equation 5.8:

\[
M = \exp\left(\int_{E_{\text{min}}}^{E(a)} \frac{\lambda \beta E}{2\pi\epsilon_0 E^2} dE\right)
\]  
(5.9)

\[
= \exp\left[\frac{\beta \lambda}{2\pi\epsilon_0} \ln\left(\frac{\lambda}{2\pi\epsilon_0 a E_{\text{min}}}\right)\right].
\]  
(5.10)

We can derive a relation between $\beta$ and the average potential difference required to produce an electron. The potential difference between the surface of the wire ($r = a$) and the point.
The track was not always on the wire, and the drift path to the wire was not always parallel to the gas during the track drift. There is therefore at least as large a spread as expected.

The proportional avalanche has a lateral extent that is at least as large as this spread. In the avalanche process itself there is a lateral development associated with the multiplication of charges, and this is mainly due to diffusion of electrons, the electrostatic repulsion of charges, and the propagation of ionizing photons. The intrinsic lateral size of an avalanche therefore depends on the gas (collision cross-section and UV absorption), on the number of charges in the avalanche and their density and on the electron energy that is obtained in the various parts of the multiplication process. If one wants to know whether the charges go fully around the wire or stay on one side, then the wire diameter also plays an important role.

The spread of avalanches on a wire has been studied experimentally by Okuno et al. [OKU79] by observing the signals from the positive ions in a segmented cathode tube surrounding the wire. In a mixture of Ar(90%) + CH₄(10%), the avalanches, created by a 55Fe source on a wire of 25 μm diameter, occupied only 100° in azimuth (FWHM) in the proportional region (total charge below 10⁶ electrons). When the voltage was raised, the avalanche started to surround the anode wire. Figure 4.6 shows how the angular width increased with the total charge of the avalanche. The increase came at smaller total charges when the concentration of the quenching gas was smaller. The influence of the UV photons was more pronounced at higher charge multiplication. The X-ray photon of the ⁵⁵Fe source has an energy of 5.9 keV and creates approximately 227 electrons. The finite size of the electron cloud produced by the X-ray photon can be neglected.

In a detailed Monte Carlo simulation of the scattering processes involved in the multiplication, Matoba et al. have shown how a small avalanche develops in three dimensions [MAT85]. In Fig. 4.7 we reproduce a picture of their electron density. In recent years the computational techniques for the simulation of such processes have been considerably advanced; see for example [GRO89]. We may expect more detailed insight into the dynamics of the avalanche as computer codes are developed that describe not only the various collision phenomena between the electrons, ions and gas molecules but also the important effect of the photons.

### 4.4 Amplification Factor (Gain) of the Proportional Wire

The multiplication of ionization is described by the first Townsend coefficient $a$. If multiplication occurs, the increase of the number of electrons per path $ds$ is

$$Z = \frac{V(a) - V(r_{\text{min}})}{\Delta V},$$

where electron multiplication begins is

$$V(a) - V(r_{\text{min}}) = \int_{r_{\text{min}}}^{r_{\text{max}}} a E(r) \, dr \quad (5.11)$$

$$= \frac{\lambda^2 \pi \epsilon_0}{\ln^2 \left(\frac{r_{\text{min}}}{a E_{\text{min}}}\right)} \quad (5.12)$$

$$= \frac{\lambda^2 \pi \epsilon_0}{\ln^2 \left(\frac{\lambda^2 \pi \epsilon_0}{a E_{\text{min}}}\right)}. \quad (5.13)$$

Let $\Delta V$ be the average potential required to produce an electron. The number of multiplications in the avalanche is

$$Z = \frac{V(a) - V(r_{\text{min}})}{\Delta V}, \quad (5.14)$$

Figure 5.39: Two views of a computer simulated electron avalanche around a 25 μm diameter wire [53].
and the number of electrons is
\[ M = 2^Z. \] (5.15)

Equating Equations 5.13 and 5.15 gives the relation between \( \beta \) and \( \Delta V \):
\[ \beta = \frac{\ln 2}{\Delta V}. \] (5.16)

\( E_{\text{min}} \) scales with density as \( E_{\text{min}}(\rho_0)\frac{\rho}{\rho_0} \) (since the mean free path of electrons scale with \( \rho^{-1} \)). Substituting this into Equation 5.13 and replacing \( \frac{1}{2\pi\epsilon_0} \) with \( aE(a) \), we obtain
\[ M = \exp \left[ aE(a) \beta \ln \left( \frac{E(a)}{E_{\text{min}}(\rho_0)(\rho/\rho_0)} \right) \right] \] (5.17)
\[ = \left( \frac{E(a)}{E_{\text{min}}(\rho_0)(\rho/\rho_0)} \right)^{aE(a)\beta}. \] (5.18)

\( E(a) \) can be written in terms of the uniform drift electric fields above and below the wire grid of pitch \( b \) [287]
\[ E_{\text{above}} - E_{\text{below}} = \Delta E = \frac{2\pi a}{b} E(a). \] (5.19)

Substituting Equation 5.19 into Equation 5.18 gives
\[ M = \left( \frac{\frac{b}{2\pi a} \Delta E}{E_{\text{min}}(\rho_0)(\rho/\rho_0)} \right)^{\frac{1}{2\pi} \beta b \Delta E}. \] (5.20)

We can see from this equation that consistency in the dimensions of the detector is important. In the case of an infinite wire grid with 40 \( \mu \)m wires at 2 mm pitch and with a symmetrical electric field of 8 kV/cm above and below the grid, a 5% spread in the wire diameter will give a gain resolution of \( \sigma/\mu \approx 100\% \). A 5% variation in the pitch also leads to a similar variation in gain. Similarly, small imperfections in the placement of the wire grid with respect to surrounding electrodes will distort the field leading to large spatial variations in the gain. Maintaining a constant gas density is also crucial hence the importance of a good cryostat. Calculating the partial derivative of \( M \) in Equation (5.20) with respect to \( \rho \) gives
\[ \frac{dG}{G} = -\frac{1}{2\pi} \beta b \Delta E \frac{dp}{\rho}. \] (5.21)

A small fractional change in density results in an approximately equal fractional change in gain. The calculation for a GEM is more difficult to perform given the more complicated
electric field. The gain can be calculated numerically by first obtaining a map of the electric field with a program like Maxwell®, and then propagating the electrons with the map while calculating the number of electrons according to Equation (5.3).

**Deviation from Proportionality**

Several processes can introduce deviations in proportionality or cause breakdown. If the gain is too high, the accumulation of positive ions can distort the electric field. The front end of the avalanche between the wire and the ions will feel a reduced field while the electrons in the tail will experience an enhanced field, moving the starting avalanche point farther away from the wire. This effect can lead to two different modes of behavior. If the photons can travel greater distances, multiple avalanches can develop simultaneously from a single event. The signal size is independent of the number of initial electrons. On the other hand, in the “limited streamer” mode, the UV photons ionize near their creation point. These photons also ionize the tail end of the avalanche, creating more positive ions which move the initial avalanche point farther and farther away. Once the streamer reaches a certain length the number electrons collected in an avalanche become almost independent of the initial number of charges. This occurs because the electric field becomes weaker as the avalanches are created farther away from the wire. Breakdown, or electric discharges, is also a limiting factor in electron multiplying detectors and there are different models of how this may proceed. The Townsend mechanism describes breakdown in terms of secondary electrons emitted when avalanche photons or back drifting ions strike the cathode [288]. These electrons lead to another avalanche. The first avalanche contains $MN_0$ electrons, where $N_0$ is the number of initial electrons generated by the event. The number of electrons in the second avalanche is

$$N_2 = \gamma M^2 N_0, \quad (5.22)$$

where $\gamma$ is feedback coefficient or number of feedback electrons created per avalanche electron and is usually called the second Townsend coefficient. The $i^{th}$ avalanche contains

$$N_i = (\gamma G)^{i-1} GN_0, \quad (5.23)$$
electrons. If \((\gamma G) < 1\), the total number of electrons collected is

\[
N_{tot} = \sum_{i=0}^{\infty} (\gamma G)^{i-1} G N_0
\]

(5.24)

\[
= \frac{G N_0}{1 - \gamma G}
\]

(5.25)

and if \((\gamma G) \geq 1\), then the sum diverges – this is the Townsend criterion for gas breakdown. However, the Townsend model cannot explain breakdown that occurs faster than the time it takes for the electron to drift from the cathode to the anode nor situations where the cathode is far removed from the anode. Another model involves “streamers” that form due the presence of ion space charges [288]. An avalanche produces a bulk of slow moving ions which distort the field. The field between the head of the avalanche and the anode is enhanced relative to the normal field. The field between the head of the avalanche and the ion space changes is decreased while the field behind these ions is enhanced. Photons from the main avalanche, if they are absorbed shortly after they are emitted, will generate new auxiliary avalanches in the regions of enhanced fields. This yields fast “streamers” in the front and end of the avalanche that result in breakdown.

5.5.4 Measurements with Case Detector

This experiments had several objectives. First, we wanted to establish that stable electron multiplication could be achieved with a two-phase LXe time projection chamber with fine wire electrode meshes. Secondly, we wanted to measure how the gain varied with the electric field. Finally, we wanted to combine these results with the two-phase measurements of electron and nuclear recoils with a two-PMT setup (described in the next section) to see if a charge readout system could be accomplished without decreasing discrimination. This section describes the fulfillment of these objectives.

5.5.5 Detector Description

The Xed1c (as seen in Figure 5.40) was similar the wavelength shifter detector (described in the last section) except that the top grid was an aluminum plate and the reflectors were
Figure 5.40: The Xed1c detector used to measure electron multiplication. This detector had a R9288 Hamamatsu PMT, four \( 40 \mu m \) wire grids and an aluminum plate.
made of PTFE. It had a Hamamatsu R9288 PMT in the liquid. In addition to the aluminum plate, it had four wire grids all of which contained 40 µm beryllium-copper wires at a pitch of 2 mm soldered onto Cirlex rings. The anode grid is directly above the liquid surface. The active region had a height of 0.95 cm and a diameter of 3.73 cm. The $^{57}$Co source was placed in the source tube underneath the detector.

The DAQ setup for charge and light measurements is described in Section 4.5. The Ortec 572 shaping amplifier was set to an amplification of 10× and shaping time of 6 µs.

### 5.5.6 Data

Gas gain data with $^{57}$Co was taken at five different liquid-to-gas extraction fields (between liquid surface and anode) – 5.9, 6.2, 6.4, 6.7, and 6.9 kV/cm – with the drift field (active region) at 0.5 kV/cm for all datasets. A summary of the data set settings is given in Table 5.2. The voltages were set such that the fields about the anode grid were symmetrical, making the equipotentials near the wire surface circular. This symmetry decreased the position dependence of the avalanche gain. The systematic uncertainty in the extraction fields and drift fields are 11% and 8.5%, respectively, due primarily to uncertainties in the distances between the wire grids. The electrons were pulled through the gate grid with 100%
Table 5.2: Data sets. $E_{below}$ is the field between the liquid surface and anode. $E_{above}$ is the field between the anode and top electrodes. All electric field values are in units of kV/cm.

<table>
<thead>
<tr>
<th>Name</th>
<th>$E_{below}$</th>
<th>$E_{above}$</th>
<th>$E(a)$</th>
<th>Effective Gain</th>
</tr>
</thead>
<tbody>
<tr>
<td>c_xed1c_0607061359</td>
<td>$-5.9 \pm 0.5$</td>
<td>$6.0 \pm 0.7$</td>
<td>$190.2 \pm 10.5$</td>
<td>$7.2 \pm 2.1_{sys} \pm 0.1_{stat}$</td>
</tr>
<tr>
<td>c_xed1c_0607061518</td>
<td>$-6.2 \pm 0.5$</td>
<td>$6.2 \pm 0.7$</td>
<td>$197.0 \pm 10.9$</td>
<td>$11.1 \pm 2.9_{sys} \pm 0.1_{stat}$</td>
</tr>
<tr>
<td>c_xed1c_0607061752</td>
<td>$-6.4 \pm 0.6$</td>
<td>$6.5 \pm 0.7$</td>
<td>$206.1 \pm 11.4$</td>
<td>$19.7 \pm 4.5_{sys} \pm 0.1_{stat}$</td>
</tr>
<tr>
<td>c_xed1c_0607061845</td>
<td>$-6.7 \pm 0.6$</td>
<td>$6.8 \pm 0.8$</td>
<td>$214.0 \pm 11.8$</td>
<td>$33.7 \pm 6.9_{sys} \pm 0.4_{stat}$</td>
</tr>
<tr>
<td>c_xed1c_0607061937</td>
<td>$-6.9 \pm 0.6$</td>
<td>$7.0 \pm 0.8$</td>
<td>$222.0 \pm 12.3$</td>
<td>$57.8 \pm 10.6_{sys} \pm 0.5_{stat}$</td>
</tr>
</tbody>
</table>

Figure 5.42: Histogram and fits of $S_1$, the charge ($Q$), and energy ($w_{max} \cdot (N_\gamma + N_e)$) of dataset $c_xed1c_0607061937$.

...
The peaks were fit with a double Gaussian function containing three free parameters:

\[ y(a, x_0, \sigma, x) = a \cdot e^{-\frac{1}{2} \left(\frac{x - x_0}{\sigma}\right)^2} + 0.3 \cdot a \cdot e^{-\frac{1}{2} \left(\frac{x - x_0 (136.47/122.06)}{\sigma (136.47/122.06)}\right)^2}. \] (5.26)

Two Gaussians are needed as 122 keV and 136 keV gamma distributions overlap. The ratio of 136.5 keV to 122.0 keV gamma event rates in the center of the active volume was measured to be 0.30\(\pm\)0.05 in a separate measurement with a two PMT detector (Xed1b as described in Section 5.2.1) and agrees with the predicted value. Use of this function assumes that both the fit means and sigmas scale with energy; the latter assumption is appropriate because the instrumental fluctuations, which scale with energy, dominate. The background rate was roughly 1/3 that of the \(^{57}\)Co event rate. The effect of background is minimized by implementing an energy and drift time cut. Here, energy is defined as \(E = w_{\text{max}} \cdot (N_\gamma + N_e)\) where \(w_{\text{max}}\) is the average energy to produced either an exciton or ion. As this quantity is not affected by recombination fluctuations, cutting in it removes background events that leak into the peak by fluctuations in recombination. The drift time includes only events at the bottom 2.5 mm of the active volume where the \(^{57}\)Co events dominate due to proximity to the source.

### 5.5.7 Results I. Basic Results

The peak voltages of the pulses from the shaping amplifier are proportional to the electron multiplication factor. The effective gas gain is defined as the fit mean voltage of the peak, divided by the voltage of the expected signal without gas gain: \(V = \epsilon e N_e / C_f\), where \(N_e = 4.45 \times 10^3\) is the number of electrons extracted from the event site, \(\epsilon\) is the electron extraction efficiency, \(e\) is the elementary charge, and \(C_f\) is the feedback capacitance. The voltage ratio is not equal to the actual gain due to short the integration time and the signal induction onto other electrodes [289]. Figure 5.43 is a plot of the effective gain as a function of the electric field on the wire surface. We were able to run the detector for more than a day at a maximum effective gain of \(\sim 58\). At very high unstable gains, multiple pulses could be seen separated in time by the electron drift time. This suggests that photon feedback on the cathode wires as the limitation on gain.
Figure 5.43: Gain versus the electric field on the surface of the wire. The highest measured stable gain was $57.8 \pm 10.6_{\text{sys}} \pm 0.5_{\text{sys}}$. The error bars represent the systematic errors.

### 5.5.8 Results II. Discrimination Projection

For a charge readout system to be viable for dark matter searches, the charge resolution of the readout must be comparable to that achieved by the typical PMT setup in order to have similar electron recoil discrimination power. In this section, we first decompose the various components of the charge signal variance of Xed1c and another detector by Sakurai and Ramsey [290]. Then, we make electron recoil discrimination projections to lower energy for hypothetical detectors with a PMT in the liquid of the same performance as in Xed1b and a charge readout system with lower electronics noise, higher gain and the instrumental fluctuations of the Xed1c and Sakurai detector. No projection is made for a CsI detector with no PMTs.

The charge variance in units of energy can be written as

$$
\sigma_E^2 = \alpha E + \beta E^2 + \sigma_{\text{recomb}}(E)^2 + \sigma_{\text{noise}}^2,
$$

(5.27)
Figure 5.44: The charge resolution versus gain. Also plotted are the resolution values with various components removed. The resolution value with statistical fluctuations, recombination fluctuations and noise removed is the square-root of instrumental coefficient, $\beta$, of equation 5.27. The error bars represent the systematic errors. Also shown are the $S^2$ resolution and instrumental component as seen in the two PMT Xed1b detector.

where $\alpha$ and $\beta$ are the coefficients of the statistical and instrumental fluctuation components, respectively; $\sigma^2_{\text{recomb}}$ is the energy dependent recombination fluctuation term; and $\sigma^2_{\text{noise}}$ is the contribution from the electronics noise. The instrumental coefficient is obtained by subtracting the total variance by all the other components at 122 keV.

The statistical coefficient, with terms representing the fluctuations in gain, extraction and number of initial ions, is

$$\alpha = \frac{bw_i}{\epsilon(1-r)} + \frac{1-\epsilon}{\epsilon(1-r)} \cdot w_i + Fw_i, \quad (5.28)$$

where $b = \sigma^2_A/\bar{A}^2$ is the relative variance of the gain of a single electron avalanche as given by the Polya distribution for $A$ [291, 292, 293] and has been determined in a Monte Carlo
simulation to be about 0.85 for 50 µm wires [294]; \( w_i = 15.6 \) eV is the average energy for generating an ion-electron pair in liquid xenon [295]; \( \epsilon \) is the liquid-to-gas electron extraction efficiency; \( r = 0.45 \) is the fraction of recombination at 0.5 kV/cm for 122 keV gammas (as determined by a separate measurement); and \( F \approx 0.05 \) is the Fano factor [296, 297]. The quantity \( w_i/(\epsilon(1-r)) \) is effectively the average energy required to generate an electron that reaches the gas phase. The extraction efficiency, \( \epsilon \), is the only quantity that changes with dataset and is equal to 0.62 in the third (and best resolution) dataset, giving \( \alpha = 0.057 \) keV.

The recombination fluctuation variance is

\[ \sigma^2_{\text{recomb}} = \frac{\sigma^2_r(E)}{(1-r)^2} E^2. \]  

(5.29)

At 0.5 kV/cm, \( \sigma_r = 0.071 \pm 0.002 \) for 122 keV gammas, giving \( \sigma^2_{\text{recomb}} = 94.0 \) keV². The calculation of \( \sigma_r \) at an energy is described in Section 5.6 and in [275, 267].

The electronics noise term in units of energy is

\[ \sigma^2_{\text{noise}} = \frac{w_i \sigma_{N_e}}{ Ae(1-r) } \]  

(5.30)

\[ = \frac{w_i C_f \sigma_V}{e A \epsilon (1-r)}, \]  

(5.31)

where \( \sigma_{N_e} = C_f \sigma_V/e \) is the rms of the noise in units of electrons, \( \sigma_V \) is the electronics noise and is the width of the test pulse distribution, \( e \) is the charge of an electron and \( C_f = 1.15 \) pF is the feedback capacitance of the preamplifier. Note that the electronics noise term is gain dependent whereas the statistical component is not. For the third data set of Xed1c, \( \sigma_{\text{noise}} = 6.11 \) keV (2640 electrons). Subtracting from \( \sigma^2_E \) the terms for statistical, recombination and noise fluctuations, we obtain for Xed1c an instrumental coefficient \( \beta = 0.0225 \). For comparison, the instrumental coefficient of the proportional light signal of the two-PMT detector described in the next section is \(~0.004\). Figure 5.44 is a plot the resolution versus gain, including resolution values with various components removed. We observe the instrumental coefficient changing with gain, which could be an indication that the background events are affecting the Gaussian fits. Also included for comparison are
Figure 5.45: The relative variances of all components (except recombination fluctuations) that contribute to the charge resolution of the data set with gain 19.7. We see that the electronics noise and instrumental fluctuations dominate at low energies. Fortunately, these components can be reduced.

the $S_2$ resolution and its instrumental component as seen in the two-PMT Xed1b detector. The great difference in instrumental components are unsurprising given the exponential dependency of the multiplication to wire diameter and pitch (Equation 5.20).

Figure 5.45 shows the relative variances of the various components (except for recombination fluctuations) as a function of energy for Xed1c at gain 20. Here, we see that contributions from gain, number of initial ions and extraction are negligible compared to instrumental and electronics noise. Fortunately, these two important components are also not intrinsic. The extraction fluctuations can be omitted by establishing a high enough electric field at the liquid surface so that extraction is 100%; this would entail thickening the anode wires to lower the gain in order to operate stably.

The same decomposition is performed on data taken with a single wire proportional counter detector by Sakurai and Ramsey. This detector was used to measure the wire
Figure 5.46: Charge resolution versus energy for detectors of various instrumental components, electronics noise and gain.

diameter dependence of the charge signal resolution using pure and CH$_4$-doped xenon at room temperature [290]. The gas was irradiated with 22 keV gammas from $^{109}$Cd. With 50 µm wires (the diameter closest to ours), an energy resolution of 3.2% was observed, corresponding to an energy variance of 0.49 keV$^2$. For the Sakurai detector, the statistical fluctuation coefficient is

$$\alpha = (b + F) \cdot w_{i}^{gas}, \quad (5.32)$$

where $w_{i}^{gas} = 0.0215$ keV is the average energy to create an ion-electron pair in xenon gas [247], giving $\alpha = 0.0193$ keV, which is lower than that of Xed1c because of the absence of fluctuations in liquid-to-gas electron extraction. Recombination fluctuations are minor in the gas phase and the electronics noise is not known – for simplicity both the recombination and noise variance terms are set to zero. This gives an instrumental coefficient of $\beta = 1.35 \times 10^{-4}$. This value is significantly smaller than that of Xed1c, which reflects the great differences in instrumental complexities.

To project the discrimination for other detectors, the electron and nuclear recoil band
data described in Section 5.2 is needed for two reasons: to obtain the energy dependent separation between the electron and nuclear recoils in $\log_{10}(S2/S1)$ and to get the intrinsic fluctuations (initial ions and recombination) in the charge and light signals. The decomposition of the signals is described in another Ph.D. thesis [60]). To create the fake data, we first simulated the intrinsic fluctuations in charge and light and then include statistical and instrumental fluctuations. We projected the performance of four detectors which have no fluctuations in extraction ($\epsilon = 1$) and drift field of 0.88 kV/cm, the field of the discrimination dataset that gives the best background rejection. We used the instrumental fluctuations of Xed1c and the Sakurai detector and varied the electronics noise and gain. Members of the Cryogenic Dark Matter Search (CDMS) collaboration have constructed charge amplifiers with a rms noise of 120 electrons by moving the primary stages of electronics into the cryostat [298]. As the electronics noise can probably be lowered to $\sim$50 electrons [299], we used this as a lower limit and the noise observed in Xed1c as an upper bound. We limited the maximum gain to 100 as groups have only been able to operate
stably with gain of several scores. Two sets of data were simulated for each detector. The first set contains electron recoil events of discrete energies between 2 and 100 keVr and was used for obtaining the charge resolution as a function of energy. The second set contains nuclear and electron recoil events (200,000 of each) simulated across a continuous band of energy between 2 and 50 keVr and is used for projecting discrimination. The discriminant used is $\log_{10}(Q/S1)$, where $Q$ is the charge. The method for calculating the discrimination efficiency is the same as that used for real data as described in Section 5.2. Figure 5.46 shows the charge resolution of the charge signal versus recoil energy at 0.88 kV/cm. The legend indicates the instrumental coefficient, noise and gain of the detectors simulated. Figure 5.47 shows electron recoil leakage versus recoil energy. Reducing the noise component (Equation 5.30) has the greatest impact on the performance. The Xed1c and Sakurai instrumental fluctuations at 5 keVee are $\sigma_{instr} = \sqrt{3E} = 0.06$ keVee and 0.75 keVee, respectively. Lowering the noise to 50 electrons and increasing the gain to 100, lowers $\sigma_{noise}$ from 6.11 keVee to 0.023 keVee and thus lowering the noise and/or increasing the gain has a much greater impact on discrimination than lower instrumental fluctuations. The detector with instrumental fluctuations of the Sakurai detector, noise of 50 electrons and gain of 100 displays a performance that is slightly inferior to that of Xed1b; the average leakage below 20 keVr is about 1.5%, whereas in Xed1b it is 0.5%. The results of this detector can be viewed as an estimate of the upper limit to the performance of this technology since full detectors are likely to have greater instrumental fluctuations and electronics noise.

In summary, we have demonstrated, for the first time, charge multiplication in a two-phase xenon detector with fine wire mesh electrodes. Also important was the stability obtained for over a day. The contributions to the charge resolution were decomposed and extrapolated to lower energies to understand which factors were important. By combining the instrumental and statistical fluctuations with intrinsic fluctuations in the light and charge signals, the discrimination performance of charge gain detectors were projected. The projections indicated that the performance of this technology can nearly reach that accomplished by the typical electroluminescence charge-readout technique if electronics noise is
greatly reduced and the gain is moderately increased.

5.6 Recombination Fluctuations in Electron Recombination

5.6.1 Introduction

Fluctuations in ionization and scintillation signals of liquid xenon detectors have been observed to be much larger than predicted by Poisson statistics. The Fano factor is defined as $F = \sigma_e^2/N_e$ [300] and describes the deviation of fluctuations from Poisson statistics. Poisson statistics corresponds to $F = 1$ and for many materials, $F < 1$ is observed (germanium, for example, has $F = 0.129$ [301]) as expected since the ionization processes are not independent. In the Case xenon detectors, the observed fluctuations in the charge signal correspond to $F > 10$ whereas the theoretical value is about 0.05 [296]. This discrepancy has been observed by many groups and has been presumed to be due to fluctuations in the fraction of recombination from event to event [243, 302]. Thomas et al proposed that the basis of the recombination fluctuations is the statistical fluctuations in the ionization density along the track of the primary particle [303].

This section describes a simple model by which the recombination fluctuations manifest in the primary light and charge signal and show how measurement of the spread of these two quantities can be used to obtain the recombination fluctuation. This method is applied to a two-phase measurement of 122 keV electron recoils at many electric fields. A binomial-like behavior in recombination fluctuation is discovered. An effort to model recombination fluctuations of this experiment is described in another dissertation [60].

5.6.2 Recombination Fluctuations Model

Radiation interacts with xenon by generating excitation and ionization. A review of the basic processes is given in Section 3.4. Recall that any electron that recombines yields a photon. This mechanism leads to a 1-to-1 anticorrelation in the calibrated $S1$ and $S2$ that can be seen across electric fields (as the mean fraction of recombination is changed) and
also at a field as the fraction of recombination fluctuates as the track structures vary from
event to event. The equations describing the number of electron \( N_e \) and photons \( N_{ph} \)
emitted can be written as

\[
N_e = N_i \cdot (1 - r) + C_1, \tag{5.33}
\]

\[
N_{ph} = N_{ex} + N_i r + C_2, \tag{5.34}
\]

\[
N_{tot} = N_e + N_{ph} = N_i + N_{ex}, \tag{5.35}
\]

where \( N_i \) and \( N_{ex} \) are the number of initial ions and excitons, respectively; \( r \) is the fraction
of initial ions that recombine; and \( C_1 \) and \( C_2 \) represent the instrumental fluctuations that
average to zero. This assumes that recombination is the only process that can occur after
the initial generation of excitation and ionization that gives a correlation between \( S_1 \) and
\( S_2 \). The initial number of ions, \( N_i \), is equal to gamma energy divided by \( w_i = 15.6 \)
eV, the average energy to create an electron-ion pair \cite{295}. To obtain the fluctuation in
recombination fraction, \( \sigma_r \), we first calculate their respective error propagation equations:

\[
\sigma_{N_e}^2 = (1 - r)^2 \sigma_{N_i}^2 + N_i^2 \sigma_r^2 + \sigma_{C_1}^2 \tag{5.36}
\]

\[
\sigma_{N_{ph}}^2 = \sigma_{N_{ex}}^2 + r^2 \sigma_{N_i}^2 + N_i^2 \sigma_r^2 + \sigma_{C_2}^2 \tag{5.37}
\]

\[
\sigma_{N_{tot}}^2 = \sigma_{N_i}^2 + \sigma_{N_{ex}}^2 + \sigma_{C_1}^2 + \sigma_{C_2}^2 \tag{5.38}
\]

The difference of \( \sigma_{N_e}^2 + \sigma_{N_{ph}}^2 \) and \( \sigma_{N_{tot}}^2 \) removes the instrumental components and gives

\[
\sigma_{N_e}^2 + \sigma_{N_{ph}}^2 - \sigma_{N_{tot}}^2 = -2((1 - r)r \sigma_{N_i}^2 - N_i^2 \sigma_r^2). \tag{5.39}
\]

Replacing the term \( \sigma_{N_i}^2 \) with \( FN_i \), where \( F \) is the Fano factor for ionization and then solving
for \( \sigma_r \), we obtain

\[
\sigma_r = \frac{1}{N_i} \sqrt{(1 - r)r FN_i + \frac{1}{2}(\sigma_{N_e}^2 + \sigma_{N_{ph}}^2 - \sigma_{N_{tot}}^2)} \tag{5.40}
\]

The Fano factor is \( \sim 0.05 \) \cite{296, 297} which makes the first term of Equation 5.40 negligible.

### 5.6.3 Experimental Setup

The Xed1h detector used for this measurement is essentially the same as Xed1b (Figure 5.1,
Section 5.2.1), except for the replacement of the PMT in the gas (R6041, QE = 7%) with one
having a higher quantum efficiency (R9288, QE = 24%). As with previous measurements, the $^{57}$Co source sat directly underneath the detector and was collimated by the PTFE rod. Data with $^{57}$Co was taken at seventeen different electric fields ranging from $\sim 0$ to 4.5 kV/cm.

### 5.6.4 Analysis and Results

The $S_1$ and $S_2$ are calibrated using the single-phase $^{57}$Co data. To calculate the recombination fluctuations, we need the variances of the peak in $S_1$, $S_2$ and $N_{tot}$. Because of the presence of 136.5 keV gammas, the peaks are fit to the double-Gaussian function with three fit parameters (Equation 5.26).

Figure 5.48 is a plot of $S_1$ versus $S_2$. The quantities are calibrated with the highest field dataset where the calculated electric field has the lowest systematic error and the
photon and electron yields change the least with field, thus providing the calibration with the lowest systematic errors. The lack of curvature in the data points shows that the calibration coefficient were constant with changing electric field and supports the idea that recombination model described in Section 5.6.2 is correct. The S1 calibration is unlikely to have drifted as this quantity is not dependent on the electric field, temperature or pressure of the system, unlike that of S2. Thus the deviations of the points from the best-fit line are more likely due to drifts in the S2 calibration. The points deviate in $S_2$ from the line by at most $\sim 2\%$. The pressure fluctuated by $\sim 0.1\%$ and thus could not have been the source of this fluctuation. It was likely due to fluctuations in the liquid level as it has been observed to fluctuate by up to 4% over the span of $\sim 10$ hours. Figures 5.49 and 5.50 show the $S_1$ and $S_2$ resolutions ($\sigma/\mu$), respectively, as a function of the electric field. Figure 5.51 is a plot of the various $\sigma_N$ values that enter into $\sigma_r$ (Equation 5.40). In Figure 5.51, we see that $\sigma_r$ exhibits a binomial-like behavior with $r$ and has a peak at $r \approx 0.4$ (as do $\sigma_{ph}$ and
Figure 5.50: \( S_2 \) resolution of 122 keV gammas versus drift field. The error bars represent the systematic uncertainty.

\( \sigma_e \), which is off by \( \sim 0.1 \) from where the peak would be for a binomial distribution. The observed variances are \( \sim 100–150 \times \) that given by binomial statistics. The systematic errors in \( r \) and \( \sigma_r \) are due to several uncertainties: the drift field due to errors in the dimensions between wire grids, the ratio of 122 keV to 136 keV gammas, and the \( S_2 \) production. The uncertainty in the highest drift field value leads to an uncertainty in the calibration in light and charge. The ratio of 136.5 keV to 122 keV gammas has been measured to be \( 0.3 \pm 0.05 \) in another experiment. The long time scale \( \sim 4\% \) fluctuations in the liquid level also contribute to the uncertainty in the calibration of \( S_2 \). The errors bars in \( \sigma_r \) and \( r \) displayed in Figure 5.52 represent the total error. The systematic error in \( r \) is \( \sim 4 \times \) larger than the statistical portion.

In this section, we described the extraction of the spread in fraction of recombination of electron recoils at an energy. Another dissertation [60] describes an attempt to model recombinations fluctuations by implementing the Thomas-Imel box model [304] on simu-
Figure 5.51: Gaussian fit $\sigma$ values versus recombination fraction. The error bars represent the systematic errors. The statistical errors are $\sim 1/5$ the systematic ones.
Figure 5.52: Recombination fluctuation versus recombination fraction
lated electron and nuclear recoil ionization tracks. A probe of $\sigma_r$ at lower energies would be useful to dark matter searches. By obtaining the energy and field dependence of $\sigma_r$, the discrimination performance of future detectors can be extrapolated by including their projected statistical and instrumental fluctuations. This analysis would also allow us to determine the electric field that achieves the optimal discrimination efficiency. Note, however, that the procedure described in this section only applies to a peak of known energy. Rather than performing this procedure, the energy dependence of $\sigma_r$ at a field can be deduced from the $\log_{10}(S_2/S_1)$ by subtracting out known contributions of statistical and instrumental fluctuations and then assuming that the remaining spread is due purely to recombination fluctuations. This procedure obtains $\sigma_r$ in a less direct fashion than the method described here but it is the only realistic way of accomplishing the task. Implementation of this procedure along with projections of discrimination performance of detectors of various light collection efficiencies is also shown in [60].
Chapter 6

XENON10

6.1 Introduction

The XENON10 collaboration, which formed in 2002, included professors, students, and technicians from Columbia University, RWTH Aachen University, Brown University, University of Coimbra, Case Western University, Princeton University, Livermore National Laboratory, Rice University, Yale University, University of L’Aquila, and Gran Sasso National Laboratory. Before building a full-scale dark matter detector, the XENON10 collaboration assembled several prototypes. The largest test detector was XENON3 [305], which was a 3 kg module with full 3D position reconstruction capability. After demonstrating proof of concept and measuring the limits of this technology, the collaboration began work in December 2005 on XENON10, a 10 kg module for taking dark matter data. The detector was largely built at Nevis Laboratories of Columbia University. Once basic tests were conducted, the detector was disassembled and reassembled at Gran Sasso National Laboratory in Italy, where it ran from April 2006 to August 2007. WIMP search data was taken and cross section upper limits were reported [13, 70].
6.2 Detector and Supporting Systems

The XENON10 experiment contained components that were similar to that of the Case setup but of larger scale and greater complexity. Other components like the Pb/poly shield was not needed for the prototypes but was required for XENON10 because of the need for passive background removal. The data acquisition was considerably more complicated as many more PMT signals needed to be recorded. The gas system of XENON10 did not differ considerable from the Case system because the storage and flow requirements were similar.

This section describes the XENON10 detector, cryostat, data acquisition system, gas system and radiation shield.

6.2.1 XENON10 Detector

The XENON10 detector was a two-phase time projection chamber with an active LXe mass (volume) of 13.5 kg (5 liters), four electrode meshes and two PMT arrays containing 89 Hamamatsu R8520 PMTs. The total xenon mass was \(\sim 22\) kg. Figure 6.1 is a 2D CAD diagrams of the detector. Figure 6.2 shows several members of the XENON10 collaboration assembling the detector in a cleanroom.

The PMTs have a 1\(''\) by 1\(''\) window covered with bialkali Rb-Cs-Sn photocathode and have a QE of 23.7\(\pm\)1.9\% at 178 nm. The effective quantum efficiency of the PMT is the QE multiplied by the electron collection efficiency of the first dynode, which is 75\(\pm\)5\%. Figure 6.3(a) shows the bottom PMT array which contained 41 PMTs and sat 1 cm below the cathode mesh. Figure 6.3(b) shows the top PMT array which contained 48 PMTs. The gap distance between the PMT windows and top mesh was 1.3 cm. This array contained one dead PMT (#48) which was not a significant lost as this PMT sat on the periphery. The high voltages were provided with a CAEN A1733. Multi-pin Kyocera feedthroughs were used to bring high voltage into the detector. Standard single-stranded kapton coated wires (from MDC) were used to bring the high voltages from the feedthroughs to the PMTs.

The reflective PTFE tube (seen in Figure 6.3(b)) had an inner diameter of 20 cm. This piece defined the diameter of the active volume. Another PTFE ring wrapped around this
Figure 6.1: Schematic of XENON10 detector and cryostat.
Figure 6.2: Assembly of XENON10 detector in cleanroom by several members of the XENON collaboration. From right to left: K. Giboni (Columbia), R. Gomez (Rice), and M. Yamashita (Columbia).
Figure 6.3: View of the bottom (top figure) and top (bottom figure) PMT arrays. The bottom and top arrays contain 41 and 48 Hamamatsu R8520 PMTs, respectively. Also seen in the bottom figure are the stainless steel meshes and PTFE reflector ring.
Figure 6.4: View of the top three of the electrode meshes and the outer PTFE ring.

The detector had four 0.203 mm thick 304 stainless steel meshes which were electro-formed and electropolished to minimize electrical discharges. The gate and cathode meshes, separated by 15 cm, defined the vertical bounds of the active volume. The anode and top meshes sat above the liquid level. The gate–anode and anode-top mesh gap distances were both 5 mm. The dimensions of mesh holes were 2.0 mm by 2.0 mm and the bars had a width of 0.182 mm. The meshes had a normal optical transparency of ~84%. The placement of the top three rings can be see in Figure 6.4.

Embedded in the inner PTFE tube were copper field-shaping rings that established a uniform field in the active region (see Figure 6.5). These were needed as the height to diameter ratio was quite large unlike that of the Case prototypes. The voltage divider
network consists of nineteen 1 GΩ resistors. Figure 6.6 shows the electric field as simulated with Comsol Multiphysics modeling software. The electric field was uniform up to $\sim 3$ mm from the PTFE walls.

The high voltages for the gate, anode, top were provided by a CAEN A1733. SHV cables carried the voltages to the cryostat and in the vacuum, PTFE insulated wire was used. In the detector, bare wire was used, which is acceptable as the LXe is an excellent insulator. A Heinzinger PNC high voltage power supply was used to bias the cathode mesh. The high voltage was brought to the detector with a custom-made PTFE insulated feedthrough.

The detector was kept at 177 K which gave a gas pressure of 2.11 atm. The liquid level was measured to a precision of 1 mm with a cylindrical capacitor and was usually kept at 2 mm above the gate electrode. For the data described in this chapter, the top and gate meshes were biased to -1.15 kV. The anode and cathode were biased to +3.15 kV and -12 kV, respectively. These voltages established a drift field of $\sim 0.73$ kV/cm and a gas field of $\sim 13$ kV/cm.
Figure 6.6: Simulated map of the electric field in XENON10. Field shaping rings keep the fields in the active region uniform. The field in the region between the cathode and bottom field array is reversed and very strong (∼9 kV/cm). This map was generated using COMSOL Multiphysics software.
6.2.2 Cryostat

A custom-made vacuum cryostat thermally insulated the detector can. The cryostat was cooled with a pulse tube refrigerator (PTR), which has a cooling capacity of about 100 W. The temperature was stabilized to $\sim 0.05$ K by a LakeShore 304 controller, several resistive Pt temperature sensors and several heaters. Figure 6.7 shows the pressure over a period of several days.

There were several heat loads in the detector: radiation, conduction, convection, xenon recirculation and resistive (ohmic). The black-body radiative load follows Stefan-Boltzmann law (Equation 4.2). The inside surfaces of the detector were primarily stainless steel which has an emissivity of $\sim 0.15$. The detector temperature was 177 K while the ambient temperature was $\sim 283$ K and the inner surface area was $\sim 0.5$ m$^2$ giving a radiative load of $\sim 23$ W. A 10-layer aluminized-mylar (same material as that used in the Case cryostat) was later added which decreased this load by roughly a factor of 10. The conductive heat load included leakages through the stainless steel cylindrical ports, electric lines, and xenon
plumbing. This heat load follows equation Equation 4.1 and was <10 W. The residue gas 
\( P \sim 2.7 \times 10^{-3} \) torr added a heat load of \( \sim 50 \) W, although the 10-layer aluminized in-
sulation likely reduced this as it likely decreased the mean-free-path of the gas molecules.
The liquification of the xenon presented another heat load which is

\[
P = \dot{M}(C\Delta T + L)
\]

where \( \dot{M} = \rho \times \dot{V} = (5.4 \text{ g/l}) \times (2.3 \text{ l/min}) = 0.2 \text{ g/sec} \) is the mass flow, \( C = 0.158 \text{ J g}^{-1} \text{ K}^{-1} \) is the heat capacity \([306]\), \( L = 95.7401 \text{ J g}^{-1} \) is the latent heat of vaporization \([306]\) and \( \Delta T = 300 - 177 \text{ K} = 123 \text{ K} \) is the temperature difference. The recirculation heat
load was about 24 W.

The PMTs introduced only a small amount of heat. The PMT bases had a voltage of
\( \sim 780 \text{ V} \) across 12.6 M\( \Omega \) voltage dividing network. The power dissipated by the PMTs was

\[
\# \text{ of PMTs} \times \frac{V^2}{R} = 89 \times \frac{780^2}{12.6 \times 10^6} = 4.3 \text{ W}
\]

The voltage divider chain for the field shaping rings emited <20 mW.

### 6.2.3 Emergency Precautions

The accidental lost of cooling power can be disastrous in liquid xenon detectors. Each
Hamamatsu 8520 PMT cost \( \sim \$1,000 \) and require many months to remanufacture. Further-
more, the evaporation of liquid xenon can lead to an explosion, causing significant damage
to the lab and resulting in the lost of xenon. Several precautions were taken to prevent this
from occurring. An emergency cooling system was built into the cryostat which consisted
of a dewar containing liquid nitrogen, a solenoid valve, and an insulated line for carrying
the nitrogen to copper coils on the detector can. The solenoid valves automatically turn on
upon loss of power to keep the detector can cooled. In addition, the detector had a burst
disc with a rupture pressure of \( \sim 4 \text{ atm} \) to protect the PMTs which implode at \( \sim 7 \text{ atms} \).
Furthermore, system vitals were constantly monitored above ground at all hours. Finally,
a text messaging system was set up to alert cell phones of system abnormalities.
6.2.4 Gas System

The XENON10 gas system consisted of a bottle rack containing two gas bottles and a gas panel containing a sealed diaphragm pump, SAES PS4 getter, mass flow controller and many high purity valves. It is very similar to the one being used at Case (see Section 4.2). Oxisorb purifiers cannot be used because of residual $^{238}$U and $^{232}$Th [307]. The amount of xenon in each bottle was tracked with a strain gauge. Figure 6.8 is a schematic diagram showing all the components and connections. Figure 6.9 is a photo of the gas panel mounted onto the lead shield.

A small dedicated LXe drift chamber for continuous monitoring of the electron lifetime was attached to the bottom of the detector. The device consisted of a solid cathode and anode, two meshes in between for shielding the solid electrodes and a series of field shaping rings. See Figure 6.10 for a picture of this device. The electrons are generated by pulsing the CsI on the cathode with a burst of light from an external xenon lamp. The lifetime value
Figure 6.9: XENON10 gas system panel, which was largely developed by members at Coimbra University.
Figure 6.10: XENON10 purity monitor developed by group at University of L'Aquila.
6.2 Light detection simulations

XENON's detector's performance was simulated with a GEANT program for its direct light and proportional light detection with a detailed detector configuration implemented (Figure 4.7). The simulation of light detection for proportional light also produced a lookup table for XY position reconstruction as in Section 6.4.1. The assumptions for the simulation are listed in Table 6.2.

is extracted from the charge signal of the anode and cathode. The device was extensively tested but was ultimately not implemented due to the head load it introduced. For more information on this device and its measurements, see [59]. See Section 6.4.1 for information on the results of the electron lifetime measurements.

6.2.5 DAQ

This section briefly describes the data acquisition system (DAQ) shown in Figure 6.11. Each PMT signal was amplified by a Phillips 776 10x amplifier (each module contained 16 inputs) which has two outputs for each input. One output was routed to one of the inputs of the SIS-3301 14-bit 105 MHz digitizer and the other was directed to the discriminator. These digitizers sampled quickly enough to capture the shape of the S1 and S2 pulses, which have widths of several tens and several hundreds of nanoseconds, respectively. Figure 6.12
Figure 6.12: The PMT traces of a nuclear recoil event. The left inset shows the individual traces of the S1 pulse which contains \(\sim 19\) phe. The right inset shows that of the single electron S2. The large pulse on the right is the S2 of \(\sim 1630\) phe. Plot courtesy of P. Sorensen [54].

shows the traces of a nuclear recoil event. Although, more information on the S1 could have been saved with faster cards (like the 500 MHz Acqiris cards used with the Case detectors) this would have cost significantly more.

The digitizer captured the 80 \(\mu\)s (maximum drift time) segment preceding the trigger and 80 \(\mu\)s segment following the trigger for each PMT. The pre-trigger section was required to record lower energy events that were triggered by the S2. The storage size of an event was 14 bits \(\times\) 88 channels \(\times\) 160 \(\mu\)s \(\times\) 105 MHz = 20,697,600 bits = 2.47 MB. Storing a 400,000 event dataset would have required nearly a terabyte of disk space. At the time of the experiment, redundant storage of this amount would have cost several thousands of
dollars. Fortunately, much of the trace does not contain useful information. A compression technique was applied (different from that used for the Case data; see Section 4.6.1) which reduced the size of the files by more than an order of magnitude. First, the segments of the signal that deviated little from baseline were set to baseline. Next the files were compressed with, gzip, which worked well with on these files as they contained much redundancy. More information on the compression procedure can be found in [60, 54]. The compression applied on Case data is described in Section 4.6.1.

Several trigger setups were explored. The S1-based trigger required a signal coincidence of $\geq 6$ on the PMTs of the bottom array. The trigger for each PMT was adjusted so that there was a $\sim 80\%$ acceptance for a single photo-electron. One problem with the S1-based trigger was the high rate of S1 only events due to a large amount of xenon outside of the active region with no electric field. Also explored was an S2-based trigger which used the sum the signal of the 34 central PMTs of the top array. This trigger system was capable of triggering off of single electrons with an efficiency of $> 80\%$ and provided a much lower energy threshold than the S1 trigger as the S2 generally contains much more light. The threshold was set to trigger off of S2’s of 4 electrons with a 100% efficiency. Because it afforded a lower threshold and avoided the S1 only events, this trigger setup was used for the WIMP search data. A more complete description can be found in [308].

6.2.6 Radiation Background Sources and XENON10 Shield

The rate of WIMP interactions is much lower than that of ambient background radiation. Most of the background in dark matter experiments is from Compton scattering of high energy gammas. Although electron recoils can be discriminated against in two-phase LXe time projection chambers, the rejection efficiency is not high enough to adequately suppress them to below expected WIMP rates and thus passive shielding is required. Passive shielding is also needed to block high energy neutrons from muon spallation and decay of radioisotopes as their response is indistinguishable from that of WIMPs. These external backgrounds are generally suppressed by running the detector underground in a plastic and lead shield.
Figure 3.2: Schematic of the LNGS underground laboratory. Entrance is via auto tunnel from the Teramo side. XENON10 is installed in an alcove of the interferometer tunnel, which presented some engineering challenges during installation due to the low ceiling height. The remote location later proved fortuitous because of the continuous construction in the main halls during the XENON10 calibration and science data-taking.

In this section, we give a general overview of background sources, a description of the XENON10 shield, a survey of background in XENON10 and the performance of the shield in removing these sources.

General Overview of Background Sources

There are many sources of background radiation. In describing the background, we divide the sources into two groups: those induced by muons and those induced by non-muon sources, primarily those of primordial heavy isotopes. An complete survey of background sources relevant to low-background detectors is given in [55].

Much tertiary radioactivity stems from three primordial sources: $^{238}$U, $^{232}$Th and $^{40}$K. The half-life of these elements are 4.57, 14.1, and 1.25 billion years, respectively. The
average concentration of $^{238}\text{U}/^{232}\text{Th}/^{40}\text{K}$ in the upper crust is 36/44/850 Bq/kg [55]. The concentration of these elements are likely to be the same in buildings and common detector materials as they are not separated in the extraction of materials from ores. A key exception is copper which is purified of these contaminants in repeated electrolytic dissolution and redeposition cycles as it has a high redox potential relative to that of U, Th and K.

Figure 6.14 describes the $^{238}\text{U}$ and $^{232}\text{Th}$ decay chains, which release a slew of alphas, betas, and gammas. Also released are neutrons from $(\alpha,n)$ reactions, in particular by alphas from U and Th on lighter elements. Nearby rock emits neutrons from $(\alpha,n)$ reactions of U and Th decay with a flux of $\sim 2$ n/m$^2$/day [309].

$^{238}\text{U}$ decays into $^{222}\text{Rn}$, which is by far the strongest airborne radioactive source. The outdoor concentration of radon is typically several Bq/m$^3$ [310]. The average indoor concentration of $^{222}\text{Rn}$ is roughly 40 Bq/m$^3$ [55]. $^{222}\text{Rn}$ is not as important because of its shorter half-life of 55.6 seconds compared to 3.82 days of $^{222}\text{Rn}$. More importantly, the progenies of $^{220}\text{Rn}$ are short lived – the half-life of $^{212}\text{Pb}$ is 10.64 hours which is by far the largest half-life of isotopes in the $^{220}\text{Rn}$ chain. On the other hand, $^{222}\text{Rn}$ eventually decays to $^{210}\text{Pb}$ which has a half-life of 22.3 years and will release two betas and an alpha before transforming into the stable $^{206}\text{Pb}$. All surfaces become radioactive from $^{222}\text{Rn}$ by the plate-out effect where by its charged radioactive daughters attach themselves to aerosols and collect onto surfaces, especially statically charged surfaces of materials such as plastic and glass. The decays at the edges of the detector can be misinterpreted as nuclear recoils and charged daughters can also migrate guided by the electric field toward grids – these events can be removed by position cuts. Covering parts with foils and keeping them in a radon free environments can lessen the contamination.

$^{40}\text{K}$ decays can decay to $^{40}\text{Ar}$ by electron capture or positron emission (11.2%), the latter of which also emits a 1.46 MeV gamma ray, or to $^{40}\text{Ca}$ by beta decay (88.8%) with a 1.32 MeV end-point energy. Potassium is an essential element for several bodily functions and $^{40}\text{K}$, comprising 0.012% of all potassium, is taken up just like regular potassium ($^{39}\text{K}$). The $^{40}\text{K}$ activity in the average human is approximately several kBq. To block the transfer
Figure 6.14: $^{238}$U and $^{232}$Th decay chains.
of $^{40}\text{K}$ to the detector via bodily fluids, gloves must be worn in handling and parts must be cleaned.

$^{85}\text{Kr}$ is produced in small amounts by cosmic rays on $^{84}\text{Kr}$. The abundance of $^{85}\text{Kr}$ in the atmosphere increased dramatically with the advent of nuclear energy but has since declined sharply due to cessation of atmospheric nuclear weapons test coupled with its short 10.76 year half-life. Fission of $^{235}\text{U}$ results in two large fragments of mass number between 90 and 140 and several neutrons. About 3 $^{85}\text{Kr}$ atoms produced per 1000 fissions. Nuclear reprocessing (to retrieve usable sources) releases $^{85}\text{Kr}$ into the atmosphere and results in an average concentration of 30 mBq/m$^3$ air [55].

Many plastic materials consist primarily of C, H and O and thus have $^3\text{H}$ (tritium) and $^{14}\text{C}$, which are radioactive. Although, plastics usually derive from petroleum sources that contain very little of these two sources, they may be introduced by additives and catalysts and by the machining of these materials [55].

The lead used as shields typically has $^{210}\text{Pb}$ which beta-decays with a half-life of 22 years to $^{210}\text{Bi}$, which in turn beta decays with a half-life of 5 days to $^{210}\text{Po}$ (see Figure 6.14). The betas of $^{210}\text{Pb}$ are unlikely to escape the lead but the 1.16 MeV beta produces a bremsstrahlung continuum with a maximum at 170 keV and characteristic x-rays with energies 72.8, 75.0, 84.9 and 87.4 keV. Concentrations of $^{210}\text{Pb}$ have been measured from detection limit up to 2500 Bq/kg [311]. Lead can be further contaminated with $^{210}\text{Pb}$ by the antimony added for hardening or by the coal used in smelting. Lead commonly used for underground experiments have a contamination of about 200 Bq/kg. Lead with activities of 5, 20 and 50 Bq/kg can be obtained at higher prices. Supplies of low activity lead produced several half-lives ago are limited to sources such as old lead pipes and ship ballast. Ancient Roman lead found in the Mediterranean sea have a $^{210}\text{Pb}$ activity of several mBq/kg, the lowest ever reported [312]. As previously mentioned, silver-tin solder can be used in place of common lead-tin solder which can have high concentrations on $^{210}\text{Pb}$.

Cosmic rays generate a slew of particles. The relative intensity of $\pi^\pm:p:e^-:n:\mu^\pm$ is 1:13:340:480:1420 at the surface of the earth [55]. The cosmic-ray-generated neutrons do
not reach underground event sites as they have a mean length of 200 g cm$^{-2}$, while muons are attenuated with 2 kg cm$^{-2}$, thus muons are the only cosmic ray daughters that affect underground experiments. Figure 6.15 shows the intensity of cosmic ray daughters as a function of depth. At sea level, the flux of muons is roughly 0.02 cm$^{-2}$s$^{-1}$ and depends on such variables as the geomagnetic latitude and the phase of the solar cycle. Figure 6.16 shows the muon intensity as a function of depth.

**XENON10 Shield**

Much of the work on the XENON10 shield was done by the Brown University group. Figure 6.17 is a 3D CAD diagram of the XENON10 shield with the sliding door open. Neutrons were moderated by a 20 cm thick layer of high density polyethylene (HDPE) weighing 1.56 tonnes. Gammas were blocked by 20 cm layer of Pb weighing 33 tonnes which surrounds the polyethylene. The standard outer lead has an activity of 560 Bq/kg. The inside layer of lead has an activity of 17 Bq/kg. The shield was supported with steel beams. One
to flat overburden by the experimental measurements of the total muon intensity. This definition and these intensities are used hereafter.

Figure 6.16: The total muon flux in depth of the various underground sites indicated [56]. The smooth curve is our global fit function to those data taken from sites with flat overburden (equation (4)).

**Table I**: Summary of the total muon flux measured at the underground sites and the equivalent vertical depth relative to a flat overburden. The horizontal lines indicate the root-mean-square deviation amongst the residuals based upon the experimental measurements.

### B. Stopping Muon Intensity

Stopping-muons are a source of background. For example, $\mu^-$ capture on a nucleus produces neutrons and radioactive isotopes. The total stopping-muon rate has contributions particularly through internal conversion (virtual-pho-

### C. Muon Energy Spectrum and Angular Distributions

In addition to the external background, the underground detectors are full of gammas from cosmic-ray interactions in the earth’s crust. To eliminate this background, we define the equivalent depth relative to a flat overburden by the experimental measurements.
4.3.5 Cryostat, Gas Handling and Purification System

The cryogenics, the gas handling and the purification system have been designed by the Columbia group.

For XENON10 a total of about 14 liters of liquid xenon is being used. The thermal insulation of the detector cryostat was carefully designed to minimize the complexity and the cost of the cooling system. Due to the narrow temperature margin (3.7K) between the liquid and solid phase, temperature control during liquefaction is especially important. Overcooling can damage the PMTs and the wire grids. A viable alternative to liquid nitrogen cooling is represented by a Pulse Tube Refrigerator (PTR). A PTR is a cooling system, which makes use of the Stirling cycle. In a PTR the Helium gas exchanges heat with a so-called orifice tube via several phases: (i) the gas is compressed in the compressor; (ii) flows in the compressor "after-cooler", where the heat is rejected to a water cooling loop; (iii) then the gas flows through the regenerator, which is basically an economizer, conserving cooling from one cycle to the next; (iv) the gas then enters the cold-end heat exchanger where heat is added to the gas from the surroundings; (v) finally the gas enters the Pulse Tube, orifice, and reservoir. These three last components produce the phase shift of the mass flow and pressure, which is necessary for cooling. The gas goes back and forth between the hot and cold ends rather than circulating continuously around a loop. Heat is lifted against the temperature gradient and rejected at the hot-end heat exchanger.
Figure 6.18: View of the detector with the shield open.
Table 6.1: Interaction lengths of high energy neutrons in polyethylene [75].

<table>
<thead>
<tr>
<th>Energy (MeV)</th>
<th>σ (barn)</th>
<th>λ_{att} (cm)</th>
<th>exp(20 cm/λ_{att})</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.1</td>
<td>28</td>
<td>0.893</td>
<td>5.3×10^9</td>
</tr>
<tr>
<td>1</td>
<td>6</td>
<td>4.17</td>
<td>121</td>
</tr>
<tr>
<td>10</td>
<td>2.5</td>
<td>10</td>
<td>7.4</td>
</tr>
<tr>
<td>100</td>
<td>0.42</td>
<td>59.5</td>
<td>1.4</td>
</tr>
<tr>
<td>1000</td>
<td>0.28</td>
<td>89.3</td>
<td>1.25</td>
</tr>
</tbody>
</table>

**Gamma rays from external sources**  The U and Th of surround rock can give an event rate of $10^4$ counts kg$^{-1}$ keVee$^{-1}$ day$^{-1}$. The Pb of the shield results in a $10^6$ reduction of the gamma background from the rock, predicting a rate of $<0.05$ dru in the detector cavity [54]. A higher rate was observed presumably due to radio impurities in the detector.

**Neutrons from external sources**  The neutron flux in Gran Sasso National Laboratory has been simulated [57]. The neutrons from radioactive decays of nearby rock at the depth of 3000-4000 mwe is roughly three orders of magnitudes higher than that derived from cosmic rays. Figure 6.19 is a plot of the neutron flux spectrum in Hall A and C. Although the neutron production rate in Hall C rock is more than 10 times lower than that of Hall A, the total rate in Hall C is only slightly lower. It was found that the neutron flux from the rock was reduced by the concrete which could be as thick as 1 m in some sections and thus, the neutrons emitted was dominated by that from the concrete. It was also determined that the rate was dependent on its water content as water is a good neutron moderator (see Figure 6.19).

Moderating the energy of a neutron to $\sim$100 keV eliminates the possibility of it generating a perceptible response in the detector. With 20 cm of HDPE, a factor 100 (60) reduction in flux of 1 MeV (6 MeV) neutrons was expected [54]. Table 6.1 gives the attenuation lengths and factors of neutrons at various energies [75]. The 20 cm Pb further reduced the neutron flux by an order of magnitude. Monte Carlo simulations predicted a neutron rate of $<0.1$ neutron/2 month/10 kg from ambient sources (non-muon) in the recoil window of 5 to 25 keVr [308].
thing inside the hall before they eventually come into the experimental setup. Those neutrons are moderated, raising the flux in the lowest energy bin.

Within the estimated uncertainties the total flux in hall C is only slightly less than in hall A for the case of dry concrete, although the neutron production rate in hall C rock is more than 10 times lower than that of hall A; above 1MeV the fluxes in the two halls are in agreement. This is due to the concrete, which indeed reduces the neutron flux from the rock significantly so that neutrons coming into the halls are mainly those produced in the concrete layer.

Table 8 shows that the neutron flux depends on the humidity of the environment. The flux in hall A is lower if the concrete is wet than if it is dry (8% and 16% water content respectively). As mentioned in the previous section, the difference in the water content of concrete does not lead to different neutron production rates. The effect seen in the flux here is caused only by moderation. Wet concrete moderates neutrons more effectively than dry concrete due its higher hydrogen content. The fluxes obtained for dry and wet concrete here show the maximum possible variation for the water content of concrete. A more realistic variation of the water content of rock in the laboratory is prepared roughly to 10–10%

Expected.

It is expected that variations of water content contribute to the results shown in Table 8. The humidity of concrete is not controlled in hall A and thus varies with the environment.

Acknowledgments

H. W. thanks the Deutscher Akademischer Austausch Dienst (DAAD) for the financial support of her Ph.D. work.

References

The distribution of neutrons produced in [34, 35, 38, 44, 45] is relevant for studies of underground sites. Nonetheless, with the advent of powerful computing, the FLUKA simulation code can be employed to compare data to experiments. Simulation studies are performed by the Case group and Chazal and Chazal.

For each site as described in Section II and used as input for the FLUKA simulation, the data are appropriately corrected for the quenching of proton-recoil energy. The neutron fluence measured in the LVD experiment before and after corrections agrees well with the FLUKA simulation performed by Wang (FLUKA) and Chazal.

Following such corrections, the FLUKA simulation of muon-induced neutrons as measured in the LVD experiment before and after corrections to the quenching of proton-recoil energy. The measured angular and energy distributions of neutrons were observed by liquid scintillator detectors and are obtained via time-of-flight. The measured angular and energy distributions of neutrons at the various underground sites are summarized in Table VII.

The primary sources of electromagnetic radioactivity are the decay of U, Th, K, and Co in the stainless steel of the detector and cryostat and the decay of $^{85}$Kr in xenon. With the nitrogen purge, the contribution from radon is insignificant. Figure 6.22 shows radon activity as a function of time. The initial xenon used had a $^{85}$Kr contamination of $\sim 25$ ppm but was later replaced with xenon from Spectra Gas certified to contain $\sim 7$ ppb of $^{85}$Kr. The Kr removal system developed by the Case group considerably improved the purity of the xenon used for WIMP search data. This xenon would have given an event rate of $\sim 2 \times 10^{-14}$ dru as opposed to $0.43$ dru expected from the xenon with 7 ppb $^{85}$Kr [308].

Neutron energy is also relevant when considering the spin of the WIMP. Neutron emission via spontaneous fission of U and Th can coincide with the emission of several MeV of gammas and thus can be discriminated against. Table 6.3 shows the expected neutron emission from WIMP decay.

![Graph showing differential energy spectrum of muon-induced neutrons at various underground sites.](image)

**Figure 17:** The differential energy spectrum for muon-induced neutrons at various underground sites. The bin width is 50 MeV.

Table VII. Summary of the fitting parameters describing the shape of the differential energy spectrum of muon-induced neutrons for the various underground sites.

<table>
<thead>
<tr>
<th>Site</th>
<th>$&lt;E&gt;$ (MeV)</th>
<th>$a_n$</th>
<th>$a_i$</th>
<th>$a_{as}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>WIPP</td>
<td>62 MeV</td>
<td>6.86</td>
<td>2.1</td>
<td>2.971 x 10^{-13}</td>
</tr>
<tr>
<td>Soudan</td>
<td>76 MeV</td>
<td>7.333</td>
<td>2.105</td>
<td>5.35 x 10^{-15}</td>
</tr>
<tr>
<td>Kamioka</td>
<td>79 MeV</td>
<td>7.55</td>
<td>2.118</td>
<td>1.258 x 10^{-14}</td>
</tr>
<tr>
<td>Boulby</td>
<td>88 MeV</td>
<td>7.882</td>
<td>2.212</td>
<td>2.342 x 10^{-14}</td>
</tr>
<tr>
<td>Gran Sasso</td>
<td>91 MeV</td>
<td>7.828</td>
<td>2.23</td>
<td>7.505 x 10^{-15}</td>
</tr>
<tr>
<td>Sudbury</td>
<td>109 MeV</td>
<td>7.774</td>
<td>2.134</td>
<td>2.939 x 10^{-16}</td>
</tr>
</tbody>
</table>
The muon-induced neutron flux for six sites (in 10 MeV.
above 10 MeV.

TABLE V: The muon-induced neutron flux for six sites (in above 10 MeV.)

<table>
<thead>
<tr>
<th>Site</th>
<th>Neutron Flux (cm(^{-2})s(^{-1}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>Boulby</td>
<td>1.0 (\times) (10^{-11})</td>
</tr>
<tr>
<td>Gran Sasso</td>
<td>1.0 (\times) (10^{-12})</td>
</tr>
<tr>
<td>Kamioka</td>
<td>1.0 (\times) (10^{-13})</td>
</tr>
<tr>
<td>Soudan</td>
<td>1.0 (\times) (10^{-14})</td>
</tr>
<tr>
<td>Sudbury</td>
<td>1.0 (\times) (10^{-15})</td>
</tr>
<tr>
<td>WIPP</td>
<td>1.0 (\times) (10^{-16})</td>
</tr>
</tbody>
</table>

For six sites (in 10 MeV.
above 10 MeV.

6.3 Data Processing

After the data was taken, we had to reconstruct detector interactions and the discrimination power depended on how accurately we were able to accomplish this. This required correcting

The fitted functions have the same form as equation (13) but with different values for parameters which are provided in Table VI. To convert the neutron production rate to the total neutron flux, one multiplies equation (13) by the average muon path length which depends upon the detector geometry.

TABLE VI: Summary of the fitting parameters describing the muon-induced neutron production rate in common detector shielding materials.

<table>
<thead>
<tr>
<th>Material</th>
<th>(P_0)</th>
<th>(P_1)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Lead</td>
<td>((7.84 \pm 2.21) \times 10^{-8})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Polyethylene</td>
<td>((6.89 \pm 1.95) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Copper</td>
<td>((4.86 \pm 0.87) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Germanium</td>
<td>((3.35 \pm 0.87) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Polyethylene</td>
<td>((2.97 \pm 0.87) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Lead</td>
<td>((3.73 \pm 0.87) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
<tr>
<td>Copper</td>
<td>((0.87 \pm 0.87) \times 10^{-9})</td>
<td>0.86 \pm 0.05</td>
</tr>
</tbody>
</table>

For cross checking purposes, two analysis groups were formed, one using ROOT (a data analysis package created by CERN) and the other using MATLAB (a commercial package created by The MathWorks). The analysis is only common in the signal compression stage.
Table 6.2: Radioactivity of various materials of the detector. The isotope activities are given as 90% C. L. upper limits. The event rate is that of the 5.4 kg fiducial volume between 1.3 and 6.6 keVee (5 and 25 keVr). This table is from [54] from which more information on internal sources can be found.

<table>
<thead>
<tr>
<th>Material</th>
<th>Mass (kg)</th>
<th>$^{235}$U, $^{238}$U</th>
<th>$^{232}$Th</th>
<th>$^{40}$K</th>
<th>$^{60}$Co</th>
<th>Event Rate (dru)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Vacuum cryostat steel</td>
<td>110</td>
<td>29</td>
<td>4</td>
<td>6</td>
<td>25</td>
<td>0.044</td>
</tr>
<tr>
<td>Detector vessel steel</td>
<td>33</td>
<td>29</td>
<td>4</td>
<td>6</td>
<td>100</td>
<td>0.176</td>
</tr>
<tr>
<td>Teflon</td>
<td>10</td>
<td>0.4</td>
<td>0.2</td>
<td>60</td>
<td>-</td>
<td>0.001</td>
</tr>
<tr>
<td>Kyocera A479SS</td>
<td>3</td>
<td>250</td>
<td>40</td>
<td>75</td>
<td>4</td>
<td>0.022</td>
</tr>
<tr>
<td>89 PMTs (PMTs + bases)</td>
<td>2.047</td>
<td>3.87</td>
<td>1.13</td>
<td>11.2</td>
<td>0.83</td>
<td>0.103</td>
</tr>
<tr>
<td>$^{210}$Pb (17 Bq/kg)</td>
<td>6370</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>0.045</td>
</tr>
<tr>
<td>Xe (7 ppb $^{85}$Kr)</td>
<td>23</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>0.430</td>
</tr>
<tr>
<td>$^{222}$Rn (&lt;5 Bq/m3)</td>
<td>&lt;0.6</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>&lt;0.010</td>
</tr>
<tr>
<td>Total (Monte Carlo)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td>0.83</td>
<td></td>
</tr>
<tr>
<td>Total (Data)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td>1.10±0.03</td>
<td></td>
</tr>
</tbody>
</table>

Table 6.3: Calculated neutron production per month from $^{238}$U and $^{232}$Th in XENON10. Table courtesy of P. Sorensen [54].

<table>
<thead>
<tr>
<th>Component</th>
<th>$^{238}$U, $(\alpha, n)$</th>
<th>$^{232}$Th Fission</th>
<th>$^{238}$U Fission</th>
<th>Total</th>
</tr>
</thead>
<tbody>
<tr>
<td>89 PMTs</td>
<td>0.25</td>
<td>0.24</td>
<td>0.29</td>
<td>0.78</td>
</tr>
<tr>
<td>89 Bases</td>
<td>2.95</td>
<td>1.22</td>
<td>0.56</td>
<td>4.73</td>
</tr>
<tr>
<td>5 HV feed-through</td>
<td>2.61</td>
<td>0.91</td>
<td>2.12</td>
<td>5.64</td>
</tr>
<tr>
<td>143 kg stainless steel</td>
<td>2.94</td>
<td>1.75</td>
<td>11.72</td>
<td>16.41</td>
</tr>
<tr>
<td>10 kg Teflon</td>
<td>0.78</td>
<td>0.50</td>
<td>0.01</td>
<td>1.29</td>
</tr>
<tr>
<td>Total</td>
<td>9.53</td>
<td>4.62</td>
<td>14.70</td>
<td>28.85</td>
</tr>
</tbody>
</table>
level of activity from $^{194}$Hg and $^{207}$Bi; these isotopes are indicative of exposure to high-energy neutrons [80], or $>10$ MeV protons [66] and were not anticipated. As described in Sec. 3.2.5, they did not significantly impact the event rate in XENON10.

Rn exclusion

The concentration of $^{222}$Rn in the LNGS cavern was measured with a Durridge RAD-7 to be as high as 140 Bq/m$^3$ in the location of the XENON10 installation. This is about $\times 2$ higher than measured in other subterranean sites, and is somewhat surprising considering the $40 \times 10^3$ m$^3$/hr flow of external air into the LNGS underground laboratory; a partial explanation is that the XENON10 installation is in a far corner, away from the main halls and ventilation ducts. Suitable Rn-exclusion measures were taken. With the shield door closed, a latex gasket made an air-tight seal such that the internal cavity could sustain a slight over-pressure (about 1.2 atm). $N_2$ gas (boil-off from a liquid $N_2$ dewar) was vented directly into the top of the shield inner cavity via stainless steel tubing, and forced to exit through the RAD-7 monitoring system at the bottom of the inner cavity. The flow rate was about 0.5 standard liters per minute (SLPM). With this arrangement, the Rn concentration inside the shield was measured to decrease to a stable value of $<5.5 \pm 1$ Bq/m$^3$ in about 30 hours, as shown in Fig. 3.13. This is actually an upper limit set by the measurement sensitivity of the RAD-7.

![Figure 3.13: Radon activity inside the shield cavity as a function of time starting at the close of shield. Plot courtesy of S. Fiorucci.](image)

Since the author was a member of the MATLAB analysis group, all the work in this section will be described from the MATLAB analysis perspective unless indicated otherwise. The author was involved primarily in defining and testing the position corrections and cuts.

6.3.1 Pulse Identification and Event Reconstruction

The PMT signal cables, HV cables, tubing for the Xe filling and circulation, $N_2$ (purge gas), He (for the PTR) and all other electronic instrumentation cables were mounted (buried) in the movable shield door. With the exception of the 6 mm $\varnothing$ blind tube for inserting the calibration source and the 6 mm $\varnothing$ tubes for the 2 leveling feet of the cryostat, all cables and tubing were fed through a z-bend (dog-leg) in the Pb. The z-bends were designed so that at least 5 cm of Pb shielding would be seen by an incident external $\gamma$, as shown in Fig. 3.14 (left). A total of 5 z-bend feed-throughs with diameter 40 mm were required, plus a 50 mm z-bend feed-through for the cryostat turbo pump. The HV cable for the cathode volatge had a...
6.3.2 Position Reconstruction

Accurate position reconstruction is important as regions near the walls of the active volume need to be excluded. The depth is provided by the drift time. The hit-pattern of proportion light on the top PMT array is used for reconstructing \((x, y)\). The proportional light was always roughly uniform on the bottom PMT array and thus provided no position information. In the analysis, two sets of \((x, y)\) position values are calculated. The first value is calculated by the “center-of-gravity” method:

\[
X_{CG} = \sum_i X_i^{PMT} \times ss2_i, \tag{6.2}
\]

\[
Y_{CG} = \sum_i Y_i^{PMT} \times ss2_i, \tag{6.3}
\]

where \(X_i^{PMT}\) and \(Y_i^{PMT}\) are the \(x\) and \(y\) coordinates of the \(i\)-th PMT, \(ss2_i\) is the integral of the S2 in the \(i\)-th PMT signal and the summation is over only the PMTs of the top array. Although the center-of-gravity coordinates have a bias toward the center of the detector, they can be calculated very quickly. The “minimum-\(\chi^2\)” method is a more accurate method that involves comparing the \(ss2_i\) to a set of those simulated in 1 mm intervals in \(x\) and \(y\) coordinates by GEANT4, a general particle simulation package (the photon propagation Monte Carlo described earlier in Section 4.3.1 generated the same results but was not used). The \(\chi^2\) quantity,

\[
\chi^2(x, y) = \sum_i \frac{(ss2_i - s_i(x, y))^2}{\sigma_i^2}, \tag{6.4}
\]

is calculated for every simulated point, where \(s_i(x, y)\) is the \(ss2_i\) simulated at position \((x, y)\) and \(\sigma_i\) is the uncertainty in the measured and simulated signals in the \(i\)-th PMT. The uncertainties in the simulated signal is insignificant compared to that of the actual signal. The fluctuations in the PMT gain and initial number of photoelectrons contribute to fluctuations in the measured signal. At first, the coordinate corresponding to the minimum value of \(\chi^2(x, y)\) was chosen as the reconstructed position but it was found that taking a
Figure 6.23: The black dots represent the actual location of simulated 10 keVr nuclear recoil events. The red dots represent the minimum-$\chi^2$ positions. Plot courtesy of K. Ni [58].

A weighted average of the positions corresponding to the three lowest $\chi^2(x,y)$ values would give a more accurate value. Figure 6.23 compares the actual and reconstructed position of 10 keVr nuclear recoil events in XENON10. Figure 6.24 shows the resolution of position reconstruction as a function of radius.

The minimum-$\chi^2$ method is very slow requiring several milliseconds per event. The processing can be sped up by at least an order of magnitude by calculating the $\chi^2(x,y)$ for only a subset of positions around an approximate $(x,y)$. For a rough location, the center-of-gravity coordinates or the position as calculated by minimum-$\chi^2$ method using a much coarser set of $s_i(x,y)$ can be used. These ideas were tested but were not well developed in time for the final analysis.

In the ROOT analysis, another method was implement which was a neural network
FIG. 8: Expected XY position sensitivity of XENON10 detector for 10 keV nuclear recoil events. The figure above shows the reconstructed event positions (red dots), compared with the original event positions (black dots). The figure below shows the dependence of position resolution ($\sigma$) on the event locations. The resolution is best at center (1.4 mm at $r = 0$), and becomes larger near the edge (2.4 mm at $r = 9$ cm).

Figure 6.24: The spread ($\sigma_r$) in reconstructed position of 10 keV nuclear recoil events as a function of radius. The resolution worsens with radius as less light is collected. Plot courtesy of K. Ni [58].

This method was more computationally intensive but was able to reconstruct the position with $\sim 100\%$ better resolution than the minimum-$\chi^2$ technique and thus was implemented in final analysis. The neural network is superior to the minimum-$\chi^2$ technique because it is able to “learn” subtle patterns in the Monte Carlo data and better extrapolate points between the positions for which the proportional light is simulated.

### 6.3.3 Basic Cuts

In the analysis, the cuts were divided into three levels. The most basic cuts, QC0 (quality cuts, level 0), removed all events that could not be considered WIMP candidates such as noise events, events that saturate the PMTs, and multiple scatters. The second level cuts, QC1, is the fiducial volume cut which excluded the edges of the active volume which have a high background rate and greater incidence of distorted events. The QC2 level cuts, removed anomalous events that may be misinterpreted as WIMP events such as those that multiple scatter with at least once vertex in the region underneath the cathode.

This section will describe QC0 and QC1 cuts. A description of QC2 will be given in Section 6.7.2.
Figure 6.25: The x represents the $t_{\text{peak}}$ and the circle represents $t_{\text{mean}}$, the first temporal moment.

**QC0 Cuts**

The QC0 level cuts had an efficiency of $\sim 99\%$ for removing undesired events. Here, we describe all the cuts comprising QC0.

**Single S1 and Single S2** Because WIMPs are expected to scatter once, only events with a single S1 and single S2 are accepted. This criteria also removes coincidence events which can be problematic as it is unclear which S1 is associated with which S2. In most instances of multiple scattering, the S2’s are distinct and the requirement of a single S2 is sufficient to remove such events. However, if the vertices are not sufficiently separated in $z$, the S2 of a double scatter may appear as a single pulse with two peaks. The pulse finding algorithm views this as a single S2 pulse and thus such events are not rejected by the single S2 scatter cut. Such pulses can be recognized by a parameter that is the difference between the mean pulse time or the first temporal moment ($t_{\text{mean}} = (\sum_i t_i V_i) / (\sum_i V_i)$, where $t_i$ and $V_i$ are the time and signal voltage, respectively, of the $i$-th time bin) and the time of the peak, $t_{\text{peak}}$. 
Figure 6.26: This is histogram of the $t_{\text{peak}} - t_{\text{mean}}$, the time of the peak subtracted by the first temporal moment (defined in the text) of S2 pulses. The outliers represent S2 pulses that have two or more peaks that was interpreted as a single pulse in the pulse finding routines. Such multiple scatters can be removed with this parameter. The vertical lines represent the conservative cut limits.
Figure 6.27: Map of final positions of S2 photons on the top PMT array. Most of photons strike two or three PMTs directly above the point at which S2 light is generated.

Normal S2’s have a roughly Gaussian shape and $t_{\text{peak}} - t_{\text{mean}} \approx 0$ sec. For double-pulsed S2’s, $t_{\text{mean}}$ lies between the two peaks while the $t_{\text{peak}}$’s will be in the middle of the larger peak. Figure 6.25 shows an S2 with two peaks with the $t_{\text{mean}}$ and $t_{\text{peak}}$ indicated. Figure 6.26 is a histogram of the distribution the $t_{\text{peak}} - t_{\text{mean}}$ distribution. Virtually all of the events that lie outside of the $\pm 0.5 \mu$s bounds are double scatter events.

If a particle scatters more than once at the same $z$ by are sufficiently separted in $(x, y)$, then the summed S2 pulse will appear as a single pulse. For such events, the hit-pattern of the top PMT array must be examined. According to photon propagation simulations, the proportional light from a single scatter is focused on the few PMTs directly above the point of generation. Figure 6.27 shows the distribution of the final photon positions
Figure 6.28: Histogram of the $\Delta r$ between the PMT receiving the most and second-most light from Monte Carlo, AmBe and $^{137}$Cs data for events with S2 between 500 and 3500 phe (40 and 120 keVr).
on the top PMT array for a single S2 pulse. On the other hand, an event with multiple scatters in the same z will generate a hit-pattern with multiple “hot spots.” A parameter which can remove such events and can be calculated quickly is the distance between the PMTs receiving the most and second-most light. The parameter, which we call $\Delta_r$, will almost always be no greater than the distance between adjacent PMTs. Figure 6.28 shows histograms of this parameter for Monte Carlo, AmBe and $^{137}$Cs data for events with a number of S2 photoelectrons between 500 and 3500, corresponding to an energy range of between roughly 20 and 140 keVr. The distance between adjacent PMTs is at most 30 cm. Any event with a high $\Delta_r$ is likely to be abnormal. The Monte Carlo showed that 0 out of $10^5$ events had $\Delta_r > 40$ cm, suggesting that actual events with $\Delta_r > 40$ cm were likely a double scatter or some other malformed event. Several events, possibly multiple scatters or other anomalous interactions, in the AmBe and $^{137}$Cs are removed by this cut. Clearly, this cut needs greater sophistication to account for multiple S2’s of different sizes. Unfortunately, this cut was developed too late to be included in the final analysis.

**Good baselines** This cut removed events whose baselines may have been calculated inaccurately by the presence of spurious pulses in first 50 samples. If samples in this segment exceeded a pre-defined threshold, the event was rejected.

**Trace Edges** This cut eliminated events with an S1 or S2 at the edges of the traces. Such pulses may have distorted pulse shape quantities as a portion of the pulse’s waveform may have been truncated.

**S2 fraction of trace area** For normal events, the S2 should comprise nearly all of the total trace area. Some events have traces with slowly drifting baselines which can result in incorrect integral and height values. This cut removes all events with $S2$ to total trace integral ratio below 0.5 or greater than 0.999.

**S2 top versus bottom** The actual and Monte Carlo data shows that~60% of the S2 signal should be collected by the top PMT array and that any large deviation is an indication
Figure 6.29: Signal trace of an event scattering immediately underneath the liquid level surface, resulting in a large $S_1$ immediately preceding the $S_2$. Such events are removed with the height/integral cut. Plot courtesy of P. Sorensen [54].

of something wrong. Events with an $S_{2top}/S_{2bottom} > 1.5$ are excluded by this cut.

**S2 width** The S2 width is determined by the distance that the electrons drift in the gas. The S2 pulse width must be between 0.45 and 2.25 $\mu$s. The range is very large because of the warping of the meshes (see Section 6.4.3).

**PMT Saturation** No WIMP event in the energy region of interest (5-25 keVr) are expected to saturate the digitizers. If any PMT signal of an event exceeds the limit of the DAQ channel, it is rejected.

**S2 height/integral** The height/integral of the pulse is a measure of the inverse of the pulse width. This cut removes broad pulses and pulses with an S1 attached to the front of the S2 (see Figure 6.29). The latter are events that occur near the liquid surface.
Figure 6.30: XENON10 background rate within $r < 80$ cm as a function of depth in the WIMP energy window.

**S1 pulse shape cut**  The S1 must meet several requirements. PMTs can have sizable dark noise pulses that are indistinguishable from primary light pulses. Such pulses are filtered out by demanding that the primary light pulse be simultaneously present in several PMTs. At least 2 PMTs must have an S1 signal of $>0.35$ phe. Also, the width at 50% maximum amplitude of the S1 must be $<0.30$ µs. Finally, the S1 rise time, defined as the time difference between the 10% maximum amplitude of the rising edge to the time of the peak, must be between 10 to 50 ns.

**Position Reconstruction Cut (χ² Cut)**  This cut excludes events with a minimum-χ² $> 2$. Events with a high minimum-χ² $> 2$ value are likely to be abnormal events.
QC1 Cuts – Fiducial Volume

Events occurring at the edges of the detector should be omitted from consideration for two reasons. Firstly, the edges will have more background events than the center due to the self-shielding afforded by liquid xenon (see Figure 3.6(b)). For example, a 1 MeV gamma has an attenuation length of $\sim 6$ cm. In XENON10, the central portion had a background rate of 0.6 events/keVee/kg/day, while the edges saw 3 events/keVee/kg/day (see Figure 6.30). Secondly, electron recoil events that occur near the PTFE walls may suffer from charge loss and be misinterpreted as a nuclear recoil. Events near the PTFE walls and the electrode meshes may also experience a different electric field than those at the center.

The fiducial volume defined in the blind analysis of the WIMP search data has a mass of 5.4 kg. The fiducial volume was determined by analyzing the $^{137}$Cs data and a portion of the WIMP search data. The fiducial volume cut has a radial limit of 80 cm and a drift time bound of 15 and 65 µs. An additional cut was used to exclude events underneath PMT #48, which was dead as the position of these events were not reconstructed well. This cut removed events with coordinates inside a circular area of radius 4.7 cm centered at the coordinate of the PMT.

6.4 Basic Performance of the Detector

6.4.1 Electron Lifetime

In the XENON10 detector, the events have a drift distance of up to 15 cm. This presented a much greater xenon purity challenge than that of the Case prototypes as these detectors had a depth of only 1 cm. A 20 cm electron drift length in Case detector would only result in a maximum charge loss of $\sim 5\%$, while in XENON10, the charge loss would have been $\sim 53\%$. The electron drift time (length) of $\sim 2$ ms ($\sim 430$ cm) was attained after several months of recirculation purification. Figure 6.31 shows the evolution of the electron lifetime over a period of about a month while recirculating xenon through the getter.
6.5 Detector Response and Stability

6.5.2 S1 dependence on the Radial Position

As in the previous case, the possible S1 dependence on the radial location of the events has to be checked and a possible empirical fitting function has to be found in order to correct the data. Making use of a set of data from a $^{137}$Cs source it has been possible to perform such analysis. Applying the method explained above, events due to photoelectric events and cutting the detector in eight concentric cylindrical cortexes, 10 mm thick, (events in the range [0,20] mm are too few to give enough statistics and are thus cut out from the analysis), it is possible to have a plot similar to the one previously called “lifetime plot”. The data are selected within a fixed range of drift coordinate. The peak positions (found for each cortex) are reported in the last plot (superimposed to all the events) and then fitted with an exponential function of the same kind as the one used for the lifetime calculation: Figure 6.32 is a plot of the S1 of this peak as a function of time. The fluctuations are less than 5% at 30 keV over a period of several months [59].

$$S1(dr) = S1(0)e^{-\frac{dr}{\rho_1}}$$  \hspace{1cm} (6.6)

6.5.3 Energy Calibration and long Term Stability

It is very important to calculate the S1 corrections explained above, especially when one is interested in the exact estimation of the energy calibration for the scintillation light
Figure 6.21: X-ray peak on the edge of the Detector. This plot has been generated by applying a radial cut rejecting all the events which happened at a distance from the center of the detector less than 85 mm.

Figure 6.22: S1 position of the X-ray peak as a function of the time. S1 has fluctuations less than 5%.

6.4.3 Light Collection

Activated Xenon

A generic problem with using external gamma ray sources to probe a detector is that the lower energy gammas (e.g. 122 keV from $^{57}$Co) do not adequately penetrate the center of the fiducial volume and higher energy sources (e.g. 662 keV from $^{137}$Cs) are more likely to Compton scatter rather than photoabsorb, giving a continuum in recoil energy rather than a peak. This problem can be solved by having the xenon be not only the target but also the emitter by inducing delayed emissions of gammas with high energy inelastic nuclear recoils. The AmBe dataset contains 40 keV prompt gammas but there are only $\sim$2000 such events.

To get a higher rate without activating detector parts, a portion of xenon can be irradiated with a very hot source and then be mixed in with the xenon that is already in the detector. Prompt lines such as the 40 keV gamma cannot be used due to the delay from transferring the xenon. Two excited states with long half-lives are $^{129m}$Xe and $^{131m}$Xe,
which emit 236 keV and 164 keV gammas with half-lives of 8.9 and 11.8 days, respectively. The $^{129m}$Xe and $^{131m}$Xe can be created by neutron capture by $^{128}$Xe and $^{130}$Xe, respectively, or by inelastic scatter off of $^{129}$Xe and $^{131}$Xe, respectively. Members at Yale University irradiated a 1 kg of xenon with a $^{252}$Cf source which emitted neutrons at a rate of $4 \times 10^4$ n/s. The xenon was activation was activated at Yale University and then shipped to Gran Sasso where a portion was included with the present xenon. This procedure was performed after taking the WIMP search data. A detail description of the preparation is provided elsewhere [236].

**S1 Signal**

The average number of photoelectrons observed in the primary light of 164 keV gammas was 2.9 phe/keVee (2.5 phe/keVee in bottom, 0.4 phe/keVee in top) at 0.73 kV/cm. The equivalent value at zero field is about 5.4 phe/keV (4.6 phe/keVee in bottom and 0.8 phe/keVee in top array) which is similar to that obtained by Xed1b, the Case prototype with two PMTs. About 8.6% of the primary scintillation photons are converted into photoelectrons. The light collection in XENON10 nearly matches that of two PMT Case prototype.

Figure 6.33 shows maps of the total primary light of 164 keV gammas collected by all PMTs and that collected by the bottom and top PMT arrays separately. Maps based on the 236 keV peaks are identical. The S1 varies with position in the expected manner. On average, the bottom array collects ~6x more light due to the internal reflection at the liquid level surface. The collection of primary light increases with depth at all radii because a lower fraction of the light is absorbed by the PTFE (see Figure 6.34). The primary scintillation collection near the bottom decreases with increasing radii, also because of the greater absorption by the PTFE. However, an opposite trend is observed near the liquid level surface, contrary to that expected from the Monte Carlo studies. The reason for this is not known. These maps shows that two separate single-variable functional corrections for drift time and radius are inadequate. The use of this map in improving the resolution is described in later sections 6.6.2.
Figure 6.33: Map of the total $S1$ (6.33(a)), and $S1$ in the top (6.33(b)) and bottom (6.33(c)) arrays in drift time and radius based on the 164 keV gammas.
Figure 6.34: S1 light collection versus depth at $r = 0$ mm as measured using the 164 keV gammas. The basic trend is followed at all radii.

**S2 Signal**

Maps of the $S2$ and $S2$ width and 10% maximum amplitude were created using 40 keV and 164 keV gammas (See Figures 6.35(a) and 6.35(c)). Neither $S2$ map exhibits the fall off with radius as expected from Monte Carlo simulations. Instead they show a bump running along the 45 degree diagonal. The $S2$ width map shows a feature of the same angle. The two values actually appear to be anti-correlated suggesting a common mechanism which could be meshes having been warped by the contraction of the supporting structures. Figure 6.36 shows scatter plots and polynomial fits of $S2$ and $S2$ width at 10% versus the azimuthal angle for 40 keV gammas events between the radius of 50 and 70 mm. The regions of the map that have a higher $S2$ and smaller $S2$ width may be regions where the anode and gate mesh are closer together. Recall that the proportional scintillation is a function of the number of electrons, electric field and drift distance (see Equation 3.19). In this case, the proportional light was more sensitive to the electric field enhancement than to the decrease in drift distance, hence we observed a higher $S2$ in regions where the grids were closer together as indicated by the lower $S2$ width. The reason for the differences between the the
Figure 6.35: Maps of $S_2$ (top-left) and $S_2$ width at 10% max amplitude (top-right) of 40 keV gammas. Maps of $S_2$ (bottom-left) and $S_2$ width at 10% max amplitude (bottom-right) of 164 keV gammas. The asymmetrical pattern is likely due to a warping of the meshes.
The use of these maps to improve the \(S2\) resolution and background discrimination are described in later sections (see Section 6.6.2 and 6.7, respectively).

### 6.5 Calibration

#### 6.5.1 PMT Gain

The gain calibration of the PMTs was accomplished by measuring the single-photoelectron response to pulses of light from internal LEDs. In the Case detectors, the calibration was performed using the single photoelectrons that spontaneously emerged from the photocathode. The use of LEDs is superior as the pulse times are known which allows for a large reduction in noise events. The LED light is controlled with a pulse generator, which had a trigger signal directed to the DAQ. The LEDs were pulsed for 6 \(\mu s\) and a 50 ns wide window was integrated 4 \(\mu s\) after the end of the LED pulse. This window had a \(\sim20\%\) probability of including one photoelectron pulse and a \(\sim2\%\) of observing two or more photoelectrons. Figure 6.37 is a histogram of the single photoelectron in one of the PMTs. The PMT voltages were adjusted to achieve gains of \(\sim2 \times 10^6\). The PMTs were recalibrated every week or two to ensure that the small drifts in signal gain were accounted for. Figure 6.38 is a plot of the PMT gain as a function of time. Even with the PMT gains equalized, they did not exhibit the same level of response due to the differences in quantum efficiencies. A correction of the relative sensitivities of the PMTs is discussed in Section 6.6.1.

#### 6.5.2 \(S1\) and \(S2\) Calibration

The nuclear recoil energy axis was not directly calibrated with a monoenergetic neutron source. As with the analysis of the Case data, single-phase relative scintillation factor (between electron and nuclear recoils) was used to calibrate the nuclear recoil energy units from the 122 keV peak. Figure 6.39 shows a histogram of the \(S1\) of the \(^{57}\)Co. The 122 keV
Figure 6.36: Scatter plot of $S_2$ and $S_2$ width at 10\% versus data of 40 keV gammas between radius of 50 and 70 mm, showing the anti-correlation between the two variables.
The single photoelectron spectrum of a PMT in XENON10. The sharp red peak is the noise and the blue broad distribution is the single photoelectron distribution.

Figure 6.37: The single photoelectron spectrum of a PMT in XENON10. The sharp red peak is the noise and the blue broad distribution is the single photoelectron distribution.

Figure 6.37: The single photoelectron spectrum of a PMT in XENON10. The sharp red peak is the noise and the blue broad distribution is the single photoelectron distribution.

gammas gives a volume averaged light yield of \( L_y = 3.0 \pm 0.1_{sys} \pm 0.1_{stat} \) pe/keVee at 0.73 kV/cm after position-dependent corrections. The event energy is scaled linearly from this peak. The electron recoil energy is \( E_{ER} = S1/L_y \), where \( S1 \) is in units of photoelectrons.

As described before, the nuclear recoil energy is \( E_{NR} = E_{ER}/L_y/L_{eff} \cdot (S_{ER}/S_{NR}) \), where \( L_{eff} \) is the scintillation yield of nuclear recoils relative to that of 122 keV gammas at zero electric field and \( S_{ER} \) and \( S_{NR} \) are the field dependent scintillation yields of electron and nuclear recoils, respectively, relative to their respective zero-field values. For simplicity, we assumed a constant \( L_{eff} = 0.19 \). For more details on previous measurements of \( L_{eff} \), see Section 3.4. The \( L_{eff} \) was also extracted from the XENON10 data and is described in [30]. At a drift field of 0.73 kV/cm, \( S_{ER} = 0.54 \) and \( S_{NR} = 0.93 \). These values give a nuclear recoil light yield of 1.02 pe/keVr at 0.73 keVr. This value is was more or less field independent as the nuclear recoil light yield does not change much with field.

The measurement of \( S2 \) of single electrons provided the charge calibration. The origin of these electrons is unknown. Some possibilities include spontaneous emission from the
Figure 6.38: This is a figure of the gains of several PMTs as a function of time. The fluctuations are less than 5%. Plot courtesy of A. Ferella [59].

cathode mesh and the delayed drift of electrons due to temporary capture by impurities. Figure 6.40 is a 2D histogram of the single electrons as a function of the $S_2$ in the top and bottom PMT arrays. The mean $S_2$ light response per electron was 23.7 photoelectrons.

6.6 Signal Corrections

Instrumental fluctuations must be removed to minimize the electron recoil leakage. Here, we discuss two such corrections on the signals to remove these fluctuations. The next section on background discrimination shows how the leakage is reduced by applying these corrections.

6.6.1 PMT Relative Sensitivity Correction

The PMTs were not all uniform in response and the quantum efficiencies of each PMT at cryogenic temperature were not given. The QE values at room temperature were given but
the QE has been observed to change with temperature in other experiments [314]. Note that the PMT gain measurements provide no information on the quantum efficiencies. Although, the absolute quantum efficiencies cannot be determined, we can obtain the relative sensitivities (RS) of the PMTs. This measurement is important for improving the position reconstruction because the reconstruction assumes a uniform response. The RS among the PMTs can be obtained by examining the response of all adjacent pairs of PMTs to the same amount of light. The MATLAB analysis used the primary light of events directly beneath the mean position of two PMTs. The ROOT analysis used proportional light. With the RS of all adjacent pairs, an overall RS map can be created, which can be used to normalize the response of all PMTs. A detailed explanation of the methods and results is explained in another thesis [54].

Figure 6.39: The $S_1$ spectrum of 122 keV gammas from $^{57}$Co in XENON10.
Figure 6.40: The 2D histogram of $S_2$ from single electrons as a function of photoelectrons in the bottom array and top array. The average $S_2$ per electron is 23.7 photoelectrons. Plot courtesy of C. E. Dahl [60].

### 6.6.2 Position Correction of Signals

As shown in Section 6.4.3, the light generation and collection varies greatly with position. Unlike the case with the Xed prototypes, the position variation can be completely removed since it has full 3D position reconstruction capability. A different multi-dimensional map is used for position correction when the RS correction is also applied. In this case, the maps are of the RS corrected signal. Figure 6.41(b) shows the histogram of the total $S_1$ with and without correcting for position with the 2D map. Figure 6.41(b) displays histograms of the total $S_2$ with and without position corrections. The improvements are also apparent in a 2D histogram plot of $S_2$ versus $S_1$ (Figure 6.42).
Figure 6.41: Histograms of S1 (top) and S2 (bottom) of the 164 keV and 236 keV with and without position corrections. The 2D maps used for correction are shown in Figures 6.33(a) and 6.35(c).
Figure 6.42: Histograms of $S_2$ versus $S_1$ of original and position corrected arrays from the activated xenon data. The 164 keV and 236 keV peaks are noticeably narrower along both axes.
6.7 Background Discrimination

Between mid-November of 2006 to December 1st, about 90 live days of $^{137}\text{Cs}$ data was taken for background calibration. A high energy veto with an upper energy threshold of $\sim 20$ keV was used to save disk space. On December 1st, a data set with an americium-beryllium (AmBe) neutron source was recorded for nuclear recoil calibration. The neutrons are produced by the alphas via $^9\text{Be}(\alpha, n)^{12}\text{C}$. These external radioactive sources were inserted through a $\sim 1/4''$ hole so as to not reduce the effectiveness of the Rn purge. There were several sets of WIMP search data. The first blind WIMP search data was taken starting in early October 2006 and ending in mid-February 2007. The WIMP search data was suspended for a short period to take aforementioned $^{137}\text{Cs}$ and AmBe data.

There are two categories of electron leakage events. The first class of events are true single-scatter electron recoils that sit on the lower end tail of the $\log_{10}(S_2/S_1)$ distribution but are in the nuclear recoil window. As will be shown, such leakage can be lessened by removing the instrumental fluctuations in the signals. The second class of leakage events are anomalous events that have a fraction of their charge signal missing. This section gives measurement of both types of backgrounds in XENON10 and the methods of minimizing them.

6.7.1 Gaussian Leakage

The procedure for determining the electron recoil Gaussian leakage in XENON10 is the same as that performed on data taken with the Case prototype. (See Section 5.2.2 for more details on the analysis procedure.) Briefly, the energy dependence of the $\log_{10}(S_2/S_1)$ mean of electron and nuclear recoils are characterized by 10th-degree polynomials and then the bands are “flattened” before being fit with Gaussian functions in discrete energy bins.

Different combinations of $S_1$ and $S_2$ corrections were explored. These corrections are described in Section 6.6. The position corrections were also tested with and without the RS correction. The nuclear recoil acceptance window is between the nuclear recoil mean ($\mu$) and three sigmas below the mean ($\mu - 3\sigma$).
Figures 6.43(a) and 6.43(b) show the electron recoil band sigma and leakage, respectively, versus energy with and without corrections. The leakage value stated in the legend is that expected for the 58.6 live day exposure. For some reason, the position correction of the $S_2$ worked better when the 40 keV electron recoil map was used, rather than the map based on the 164 keV gammas. As shown in the Section 6.4.3, the two maps are different although the diagonal feature is present in both.

### 6.7.2 Non-Gaussian Leakage

**Sources of Non-Gaussian Leakage Events**

In the $^{137}$Cs data, the number of leakage events observed in the nuclear recoil acceptance window is larger than that expected from the tail of the Gaussian distribution. The $^{137}$Cs data has 2273 events between 4.5 and 26.9 keVr (the WIMP energy window) after application of the first two sets of quality cuts (QC0 and QC1, as described in Section 6.3.3). In the tail of 0.5% Gaussian leakage, 23 events are observed corresponding to a leakage of $\sim$1%. These anomalous events may be misconstructed events. As previously mentioned, an electron recoil event can appear as a nuclear recoil event if some of its charge is lost. The events that lose charge due to interactions with the edges of the detector can be removed as their position is known. Events that scatter at least once in the active volume and at least once in a region where the primary light can be collected but not the ionization (Figure 6.44). Unfortunately, such events cannot be omitted by a simple fiducial volume cut. Most of these events (which we call “gamma-x”) are those that have a vertex in the region between the cathode mesh and bottom PMT array (reverse-field region) where the electric field drifts the electrons downward. Some of these events may also be gammas that scatter once in the pocket of liquid xenon of the high voltage divider of the field shaping ring (see Figure ).

**QC2 Cuts**

Cuts were developed to these anomalous events without severely lowering the acceptance for WIMPs. To accomplish this, the cuts were applied not only on the events of the $^{137}$Cs
Figure 6.43: Top – The Gaussian width of $\log_{10}(S2/S1)$ versus energy of the $^{137}$Cs data with and without corrections. Bottom – The electron recoil leakage fraction in the 50% acceptance nuclear recoil window as a function of energy, with and without corrections to the instrumental fluctuations as described in the text. The position dependence of $S1$ was removed with the 164 keV 2D $t_{drift} - r$ map. For $S2$, the position dependence was removed with the 40 keV $x - y$ map. The PMT relative sensitivity correction was also applied on the $S2$. 
Figure 6.44: Diagram of a “gamma-x” event. Such an event scatters at least once in the region where charge cannot be collected (such as the space between the cathode mesh and bottom PMT array) and at least once in the fiducial volume. A portion of the charge does not participate in the generation of proportional light and thus the event has a lower $\log_{10}(S_2/S_1)$ than expected.

Events but also on the nuclear recoil events of the AmBe dataset (with the nuclear recoils selected by a $\mu \pm 3\sigma$ cut in $\log_{10}(S_2/S_1)$). For both data sets, only events in the tail of 0.1% Gaussian leakage were examined.

**Reverse-field region event cut (gamma-x cut)** Background simulations of gammas from PMT and other internal sources give a gamma-x event rate of $\leq 1$ mdru [308]. The expected number of observed gamma-x events is $10^{-3}$dru$\times$10keVee$\times$5.3kg$\times$59 days = 3 events. Fortunately, some of these events can be removed by examining the hit pattern on the bottom PMT array. Primary light emitted from the region between the cathode
We have identified a large number of events with secondary scatters over 2-3 PMTs near Liquid Xenon pocket in the Teflon Can - small hole made for resistor chain might be filled with liquid Xenon

and bottom PMT array is likely to be concentrated on a few nearby PMTs unlike that emitted from the active volume. Such events can be removed by cutting in a parameter which measures the localization of primary light on the bottom array. Several such parameterizations were explored. Some simpler parameters were $\text{std}(ss_{1\text{bottom}})/S_{1\text{bottom}}$ and $\text{max}(ss_{1\text{bottom}})/S_{1\text{bottom}}$, where $\text{std}(ss_{1\text{bottom}})$ is the standard deviation of $ss_{1}$ of the bottom PMT array. $\text{max}(ss_{1\text{bottom}})$ is the size of the largest $ss_{1}$ of the bottom PMT array. Both parameters increase with greater localization of light. Figure 6.46 is a scatter plot of $\text{max}(ss_{1\text{bottom}})/S_{1\text{bottom}}$ for the $^{137}\text{Cs}$ and $^{252}\text{AmBe}$ data. This plot shows that some of the leakage events (indicated by green dots) of higher values of this parameter can be eliminated with a cut (blue line) without severely reducing the nuclear recoil acceptance. These simpler cuts were not implemented in the secondary (MATLAB) blind analysis. To remove gamma-x events, the ROOT analysis group used $\text{std}(ss_{1\text{bottom}})/S_{1\text{bottom}}$ which was also
Figure 6.46: Scatter plot of $\max(ss_{1\text{bottom}})/S_{1\text{bottom}}$ of the $^{137}\text{Cs}$ (black) and AmBe (red) events (only those in the WIMP window) as a function of energy. The electron recoil leakage events of the $^{137}\text{Cs}$ dataset are indicated by green dots. The events outside the bounds of this cut (blue line) are indicated by larger dots.

A more sophisticated and efficient set of cuts for removing gamma-x events was developed by Peter Soreneon of the MATLAB analysis group. Only a brief description is given here (more details can be found in [54]). The parameters are $G_n = \sum_i ss_{1i}/S_{1\text{bottom}} \cdot \sqrt{ST}$ where the sum is over the $n$ bottom array PMTs that receive the most light. The $\sqrt{ST}$ is present to reduce the energy dependence of the parameter. For $n = 1$, the parameter is more or less the same as the parameter described in the previous paragraph. For $n \geq 2$, the standard deviation of the x and y coordinates of the $n$ PMTs must also fall within some
Figure 6.47: Plot of $G_1$ (as defined in the text) as a function of energy of the $^{137}$Cs data. The red line indicates the upper bound for 95% acceptance as determined with a Monte Carlo. The leakage events in the 0.1% tail of Gaussian leakage are indicated by the red circles, four of which are excluded by this cut. The blue x's indicate extremely non-Gaussian events that lie beyond the $\mu - 3\sigma$ lower bound of the WIMP acceptance window. The arc of events on the upper left are events with all of the S1 light concentrated on one PMT. Plot courtesy of P. Sorensen [54].

upper limit. Cuts for $n = 1, 2, 3$ and 4 were calibrated with the help of a Monte Carlo. Figure 6.47 shows $G_1$ as a function of energy of the $^{137}$Cs data. Figure 6.48 shows the $G_2$ as a function of energy and the standard deviations of the PMT coordinates.

**Partial Gamma-X Cut** There exists a region within about 1 cm from the cathode where only a small portion of the ionization can escape. Multiple scatters with a vertex in this region will yield a small S2 pulse trailing the S1 by the maximum drift time. This secondary S2 may be $< 5\%$ the size of the primary S2 and consequently be ignored by the multiple S2 cut described earlier. The partial gamma-x cut removes such events.
S1 edge-to-all ratio cut  This cut removes events that may have scattered at the edges of the detector and thus may have lost some charge. The parameter is the sum of the S1 signal in PMTs on the outer edge of the bottom array divided by the total S1. The efficiency of this cut was determined with the same Monte Carlo as that used for the reverse-field region event cut. Figure 6.49 shows the hit-pattern of an event with a large fraction of light on the outer ring of PMTs.

S1 top fraction cut  The actual data and Monte Carlo photon propagation simulations show that most of the light should fall on the bottom PMT array. Large deviations would indicate something severely incorrect with the event. One anomaly that may result in a larger than expected proportion of the primary light falling on the top array is a scattering in the gas. Such occurrences are expected to be rare as the mass of the gas is much less than that of the liquid. Figure 6.50 shows the fraction of the S1 on the top PMT as a function
The $S_1$ $E_2A$ cut is similar in principle to the $S_1$ $RFR$ cut. In this case, an index $E_{frac} = S_1/ S_1$ was defined to reflect the extent to which an individual event might have scattered at the very edge of the detector, prior to scattering in the fiducial target. As described in Sec. 4.6.2, events at the very edge of the active target are expected to have reduced efficiency for drifting ionization to the liquid-gas interface. An example hit pattern of a non-Gaussian event electron recoil event that was cut by the $E2A$ cut is shown in Fig. 4.38. Events that are cut by $E2A$ are often (but not always) also cut by $RFR$. The performance of the $E2A$ cut is shown in Fig. 4.40 for the $^{137}$Cs calibration data.

![Figure 4.38: The $S_1$ hit pattern for an event tagged by the QC2 $E2A$ cut. The anomalous concentration of signal near the edge on the bottom array is highly unlikely for a "normal" single-scatter event. The event narrowly missed the threshold for the $RFR$ cut.](image)

Plot courtesy of P. Sorensen [54].

Figure 6.49: PMT hit pattern of an event with a large portion of the primary light concentrated at the edge indicating that they may have scattered at the edge of the detector and thus may have suffered charge loss. Such events are removed with the $S_1$ edge-to-all ratio cut.

of energy of the $^{137}$Cs dataset with the boundaries of 99% cut efficiency. The previously described Monte Carlo was used to help define the bounds.

**Pulse Area Fraction**  This cut parameter is $(S_1 + S_2)/A$ where $A$ is the total integral of the trace. Ideally most of the light should be from the primary and proportional scintillation signals. Events with values much lower than 1 may have problems with noise or may include light from abnormal sources. Some extraneously sources of light are proportional scintillation from electrons spontaneously emitted from the cathode and those photoionized by the proportional scintillation. Figure 6.51 shows this parameter plotted as a function of energy.

**$S_1$ width at 10% maximum amplitude**  This cut removes events with an $S_1$ width at 10% maximum amplitude of $>150$ µs. This cut was motivated by the observation of $10 \log_{10}(S_2/S_1)$ with $S_1$ width $>200$µs. Assuming that the $S_1$ decay probability distribution follows a simple exponential of fall time $\tau = 25$ ns, the probability of at least
one photoelectron occurring $> 200$ ns after the start of the pulse of $N$ photoelectrons is
\[ \sum_i^n (1 - \exp(200/25))^N - i (\exp(200/25))^i \approx \exp(200/25) = 3.4 \times 10^{-4}. \]
About ten times more events with an S1 width $> 200 \mu$s was observed than expected and the cause is unknown.

**Nuclear recoil Acceptance of QC2 cuts**

The QC2 cuts were applied to the nuclear recoils events of the AmBe data to determine their affect on WIMP acceptance relative to the QC0 and QC1 cuts. To exclude electron recoil events, a $\mu \pm 3\sigma$ cut was applied around the nuclear recoil band. The reverse-field-region and edge-to-all cuts are responsible for most of the reduction in acceptance. Figure 6.52 is a plot of the nuclear recoil acceptance of the MATLAB QC2 cuts as a function of
Figure 6.51: Plot of the pulse area fraction parameter versus energy of the $^{137}\text{Cs}$ data. The red line indicates the lower bound. The excluded event has a large amount of after-pulsing following the S2. Plot courtesy of P. Sorenesen [54].

6.8 Dark Matter Limit

A blind analysis is important because it removes any opportunity for biasing the results. The blind portion of the WIMP search was locked away until all the cuts were defined. To check the quality of the WIMP search data without spoiling the analysis process, only sections of energy beyond the dark matter energy range were examined. By studying the $^{137}\text{Cs}$ and the unblind portion (~16 live days) of the WIMP search data, the cuts were fine tuned. Also set before the unblinding was the energy range (4.5–26.9 keVr), energy binning and the WIMP acceptance window ($\mu - 3\sigma$). Before unblinding, the ROOT analysis was deemed by vote to be the primary analysis whose results would be shown in publication.
In this section, we discuss the results of the ROOT and MATLAB blind analysis of the 58.6 live days of WIMP search data.

### 6.8.1 WIMP Candidate Events of ROOT Analysis

The ROOT blind analysis revealed 10 candidate events in the WIMP window out of a total of 1815 events between 4.5 and 26.9 keVr after applying all cuts (QC2 cuts are described in [13]). Figure 6.53 is a plot of $\Delta \log_{10}(S2/S1)$ ($\log_{10}(S2/S1)$ centered on the electron recoil band centroid) as a function of energy of the WIMP search data. All leakage events can be explained as electron recoil background or as anomalous events. Five of the events (3, 4, 5, 7, and 9), which sit just underneath the upper bounds of the window, are consistent
For a WIMP mass of 100 GeV/c², the upper limit is $8.8 \times 10^{-44}$ cm², a factor of 2.3 lower than the previously best published limit [27]. For a WIMP mass of 30 GeV/c², the limit is $4.5 \times 10^{-42}$ cm².

Figure 3 (color online). Results from 58.6 live days of WIMP search in the 5.4 kg LXe target. The WIMP signal region was defined between the two vertical lines (4.5 to 26.9 keV nuclear-recoil equivalent energy) and the two zigzag lines (about 50% nuclear-recoil acceptance).

Figure 4 shows the 90% C.L. upper limit on considering all ten observed events, with no background subtraction WIMP-search data set. To set conservative limits on the Monte Carlo predicted spectrum and the measured energy window, from the 58.6 live days of WIMP search data. (small dots) All events. (b) Events in the WIMP-signal region before the software cuts. (c) Events in the WIMP-signal region after the software cuts. The leakage events. The ones which are labeled blue (near the top blue line) are those which are consistent with the Gaussian fit. The two vertical lines represent the 4.5 to 26.9 keV energy bounds.

FIG. 3 (color online). Results from 58.6 live days of WIMP search in the 5.4 kg LXe target. The WIMP-signal region was defined between the two vertical lines (4.5 to 26.9 keV nuclear-recoil equivalent energy) and the two zigzag lines (about 50% nuclear-recoil acceptance).

6.8.2 WIMP Candidate Events of MATLAB Analysis

Only a brief summary of the blind MATLAB analysis will be given here; a complete account can be found in [54]. Figure 6.55 is a plot of $\Delta \log_{10}(S2/S1)$ versus energy of the WIMP signal events.
FIG. 2 (color online). Position distribution of events in the 4.5 to 26.9 keV nuclear-recoil energy window, from the 58.6 live days of WIMP-search data. (small dots) All events. (+) Events in the WIMP-signal region before the software cuts. (⊕) Events remaining in the WIMP-signal region after the software cuts. The solid lines indicate the fiducial volume, corresponding to a mass of 5.4 kg.
Figure 5.2: The XENON10 58.6 live day (blind) WIMP search result from the secondary (Matlab) analysis. There are 1716 single-scatter events in the energy range $4.5 - 26.9$ keVr. Seventeen WIMP candidates remained found in the acceptance box. Of these, twelve events are fully consistent with the Gaussian electron recoil background (blue circles). Event numbers 1, 2 and 5 with \( \Delta \log_{10}(S_2/S_1) < -0.6 \) (red circles) are clearly inconsistent with the electron recoil background. Event numbers 3 and 4 are very close to the 99.9% contours (blue line), and could be reasonably expected fluctuations of the electron recoil background population. The two vertical lines represent the 4.5 and 26.9 keVr energy bounds. These consistent with the Gaussian tail. Events 3 and 4 are may also be consistent with Gaussian leakage as they sit near the 99.9% contour. Events 1, 2 and 5 are very unlikely to be members of the tail. The two vertical lines represent the 4.5 and 26.9 keVr energy bounds.

The WIMP search acceptance box was drawn by making a Gaussian fit to the transformed discrimination parameter \( y' = \Delta \log_{10}(S_2/S_1) \) for the single-scatter AmBe calibration data, in discrete bins. The bin edges were set at 2, 3, 4, 5, 6, 8, 10 and 12 keVee, using 2.2 S1 phe/keVee as obtained from the \(^{137}\text{Cs}\) full-energy peak. This encloses the nuclear recoil equivalent energy range $4.5 - 26.9$ keVr. The fits to \( y = \log_{10}(S_2/S_1) \) are shown in Fig. 2.18 for the untransformed coordinates. The Gaussian fits for the transformed coordinates are not significantly different aside from a shift in the centroid. The \(^{137}\text{Cs}\) calibration data centroid values in each bin are shown in Table 5.1 for both \( y \) and \( y' \). The actual values of \( y' \) in each bin are slightly different from the expectation of \( y' = 0 \). This is an expected consequence of the finite bin width, since all data within a given bin were transformed according to the bin centroid. The ER-subtracted centroids of the WIMP
Figure 6.56: Position of events in the 4.5–26.9 keV nuclear recoil energy window from the MATLAB analysis. The blue circles indicate the WIMP candidates (also circled blue in Figure 6.55). The events with red circles are those below the 99.9% contour and are unlikely to be members of the Gaussian tail. The blue lines represent the boundaries of the fiducial volume.
search data after the application of all cuts. After applying the cuts, 1716 events remain between 4.5 and 26.9 keVr, 17 of which are in the nuclear recoil window. Twelve of the events which are near the 99.5% contour are consistent with that expected from Gaussian leakage. Events 3 and 4 may be consistent with the Gaussian tail as they are near the 99.9% contour. Events 1, 2, and 5 are highly unlikely to be members of the tail. Events 1, 2, and 4 are events that the reverse-field-region cut would have removed if the cut had slightly tighter bounds (and consequently a slightly lower nuclear recoil acceptance). Event 1 was removed in the ROOT analysis by a S1 hit pattern cut in QC2. Event 2 is also present as a WIMP candidate in the ROOT analysis (also as event #2) and, as previously stated, likely to be a reverse-field-region event as its reconstructed position is near the bottom of the fiducial volume. Event 4 was cut by QC2 of the ROOT analysis. Event 3 was removed by the QC0 of the ROOT analysis. In this event, the S1 occurred just before the S2 and, as a result, the two pulses was interpreted as a single pulse. The fake S1 pulse was an large pickup of coherent noise that should have been spotted as it consisted of 5 photoelectrons spread out among 18 PMTs, which is roughly 0.3 pe/PMT – the chances of an actual S1 giving such a pattern is small. Event 5 was admitted because its position ($r = 78$ mm) was reconstructed incorrectly due to its location near the dead PMT. The neural network algorithm of the ROOT analysis placed this event at $> 85$ mm, well outside the fiducial volume. There are more Gaussian leakage events in the MATLAB analysis because a position correction for the S2 was not included before the unblinding. This occurred because of the discrepancy between the 40 keV and 164 keV maps which lead to an uncertainty as to which would work best and, in the end, we decided to not use a correction and devote our efforts towards the removal of anomalous events. The expected number of leakage events with and without the S2 correction are $7.4^{+1.8}_{-1.5}$ and $12.1^{+1.9}_{-1.6}$, respectively [54]. After correcting the S2, only 6 Gaussian leakage events remained which is consistent with expectations.
Although we treated all ten events as WIMP candidates in calculating the limit, none of the events are likely WIMP interactions. $\Delta \log_{10}(S2/S1)$ values for 5 events (compared to $30 \text{ GeV}/c^2$).

As explained in the previous section, all 10 leakage events of the ROOT analysis could be explained without invoking the existence of WIMPs but to be conservative in the calculation of the cross section upper limit, we assumed all those events to be WIMP scatters. The upper limit was calculated using the “maximum-gap” method [315] with standard assumptions on the galactic halo (see Section 2.3.1 or [170]). The statistical methods of Feldman and Cousins were not used as the background was not measured well enough for a statistically meaningful subtraction. The maximum gap method is useful for small signal experiments containing an unknown background. It provides a way of including energy regions that are likely to be least affected by the background. Here, we give a brief explanation on the application of this method. Let the energy of the $i$-th leakage event be $E_i$ ordered from lowest to highest energy. The $i$-th gap is defined as the expected number of events of the

Figure 4. Spin-independent WIMP-nucleon cross section upper limits (90% C.L.) versus WIMP mass. Curves are shown for the previous best published limit [upper (blue)] [27] and the current work [lower (red)], assuming a constant 19% 

6.8.3 Calculation of WIMP cross section upper limit 

L_eff. The shaded area is for parameters in the constrained minimal supersymmetric models [6,29].
energy gap $E_i - E_{i+1}$:

$$x_i = \int_{E_{i+1}}^{E_i} \frac{dN}{dE} dE,$$  \hspace{1cm} (6.5)

where $\frac{dN}{dE}$ is the recoil count spectrum scaled by a particular cross section, $\sigma$. The maximum gap, $x$, is the largest $x_i$. A particular $\sigma$ is rejected if it gives a maximum gap that is highly unlikely. In this paper, Yellin derives the probability of a maximum gap being smaller than an observed $x$ for a total expected number of events $\mu$ as

$$C_0(x, \mu) = \sum_{k=0}^{m} \frac{(k \mu - \mu)^k e^{-kx}}{k!} \left(1 + \frac{k}{\mu - kx}\right)$$ \hspace{1cm} (6.6)

where $m$ is the greatest integer $\leq \mu/k$. The 90\% confidence upper limit for an observed $x$ is the $\mu$ ($\sigma$) that gives $C_0 = 0.9$. For the XENON10, WIMP search data, the energy bounds of the maximum-gap is 4.5 and 15.5 keVr, of events 1 and 2, respectively (as seen in Figure 6.53). Recall that WIMP-nucleon coupling is characterized by two terms: spin-independent (scalar) and spin-dependent (axial vector) (see Section 2.3.1). Figure 6.57 shows the spin-independent WIMP-nucleon cross section limit as a function of WIMP mass. Also shown is the CDMS result of 2005 which was the previous best at the time [316]. CDMS has since exceeded the results of XENON10 [14].

Figure 6.58 shows the WIMP-nucleon cross section upper limits for pure-neutron ($a_p = 0$) and pure-proton ($a_n = 0$) couplings. Both $^{129}$Xe (spin-1/2) and $^{131}$Xe (spin-3/2) contain unpaired neutrons and thus XENON10 is most sensitive to neutron spin-dependent couplings. The 58.6 live data sets new limits across all WIMP mass. The WIMP-neutron limit has a minimum of $\sim 6 \times 10^{-39}$ cm$^2$ at $\sim 30$ GeV c$^{-2}$. The largest systematic uncertainty is attributed to $L_{eff}$ and could raise the limit at 30 GeV c$^{-2}$ (100 GeV c$^{-2}$) by $\sim 15\%$ ($\sim 18\%$). Without unpaired protons, XENON10 is unable to set a new WIMP-proton limit. It exceeds the sensitivity of CDMS and ZEPLIN-II and approaches that of NAIAD, PICASSO, KIMS, and COUPP.

The data places a constraint on $a_p-a_n$ space at a certain mass. The expected number of events is a function of $a_p$ and $a_n$:

$$N = Aa_p^2 + Ba_pa_n + Ca_n^2,$$ \hspace{1cm} (6.7)
protons and neutrons, respectively.

The spin enhancement factor in this case is:

\[ a = 0.46 \]

and

\[ a = 0.34 \]

taking into account the strange quark contribution to the nucleon spin, as measured by the EMC collaboration and given in [26] for coupling to protons and neutrons, respectively.

Figure 6.58: Upper limit 90% confidence level WIMP-neutron (left) and WIMP-proton (right) cross sections as a function of WIMP mass based on the data of XENON10. The black line and dashed red line represent the limit with the Bonn A and Nijmegen II potentials, respectively. Also shown are the results from the CDMS (circles), KIMS [63] (triangles), NAID [64] (squares), PICASSO [65] (stars), COUPP [66] (plus), SuperK [67] (crosses), as well as the DAMA allowed region. The expected cross section on the first time approaching the theoretically predicted parameter space for neutralinos.

The theoretical regions (blue filled) for the dark halo parameters (green filled region) [68]. The theoretical regions (blue filled) for the dark halo parameters (green filled region) [68]. The expected cross section on the first time approaching the theoretically predicted parameter space for neutralinos.

The expected cross section on the first time approaching the theoretically predicted parameter space for neutralinos.

\[ N_{Xe} = Aa_{Xe}^2 + Ba_{Xe} + Ca_{Xe}^2 \]  

with A, B, C being constants of integration of the differential event rate \( dR/dE \) over the relevant energy region, in our case 4.5 keV–27 keV nuclear recoil energy.

Figure 2 (right) shows the allowed regions inside ellipses of \( a_p \) vs. \( a_n \) parameter space for a fixed WIMP mass. We follow [18] and express the expected number of recoil events \( N_{DP} \) as a function of \( a_p \) and \( a_n \):

\[ N_{DP} = Aa_p^2 + Ba_p + Ca_n^2 \]
neutron couplings are the most sensitive to a WIMP mass of 30 GeV/c². Using a local WIMP model and a Maxwell-Boltzmann velocity distribution, we can probe a heavy Majorana neutrino with a mass larger than 30 GeV/c² as the dark matter target.

This work was funded by NSF Grants No. PHY-03-02646 and No. PHY-04-00596, the CAREER Grant No. PHY-0542066, the DOE Grant No. DE-FG02-91ER40688, by the Swiss National Foundation SNF Grant No 20-118119, by the Volkswagen Foundation (Germany) and by the FCT Grant No POCI/FIS/60534/2004 (Portugal). We would like to thank the LNGS/INFN staff and engineers for their help with the experiment.
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