

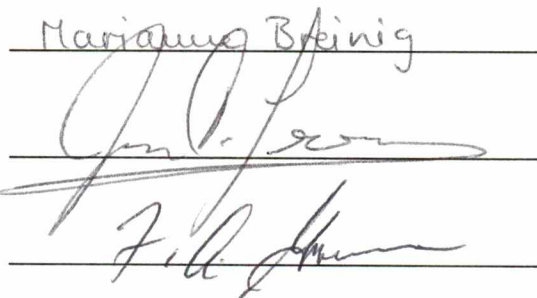
To the Graduate Council:

I am submitting herewith a dissertation written by Leticia S. Pibida entitled "Coincidence Measurements of Cascade Auger Electron Emission Processes in Photon-Atom Interactions and Autoionization of Electrons from Doubly Excited States of Helium Produced by Fast Ion Impact". I have examined the final copy of this dissertation for form and content and recommend that it be accepted in partial fulfillment of the requirements for the degree of Doctor of Philosophy, with a major in Physics.



Dr. Ivan A. Sellin, Major Professor

We have read this dissertation
and recommend its acceptance:



Accepted for the Council:



Associate Vice Chancellor
and Dean of the Graduate School

Coincidence Measurements of Cascade Auger Electron Emission Processes in
Photon-Atom Interactions and Autoionization of Electrons from Doubly Excited
States of Helium Produced by Fast Ion Impact.

A Dissertation
Presented for the
Doctor of Philosophy
Degree
The University of Tennessee, Knoxville

Leticia S. Pibida
December 1997

Dedication

I want to dedicate my dissertation to: My daughter Melissa Minniti, my parents Norma and Juan Carlos Pibida and my husband Ronaldo Minniti.

Acknowledgments

I want to especially thank my advisor, Dr. Ivan Sellin and Dr. Ralf Wehlitz for all the help and knowledge they shared with me through all these years. I gratefully acknowledge the frequent consultation of Dr. Marianne Breinig, Dr. Jon Levin, Dr. Joachim Burgdörfer and Dr. Stuart Elston, as well as their help in many aspects of my work. I thank Dr. Leif Liljeby for his contribution to the development of the data acquisition system and the other members of the UT Atomic Physics group, R. Minniti, D. Desai, J. Y. Lim and L. Straus, for their help during the experimental runs. Brad Armen provided valuable theoretical advice.

A note of appreciation for the help of many people in the physics department needs to be made specially to Mrs. Frances Cole (secretary of the Atomic Physics group), Mr. Frank Spencer (from the UT machine shop) and to all the secretaries in the physics department office.

Finally, I will also want to thank my former licenciante advisor Dr. Ignacio Nemirovsky for introducing me to the field of Atomic Physics, that I enjoy so much.

This work was supported by NSF and U.S. DOE, Office of Basic Energy Science, Division of Chemical Science, under contract number DE-AC05-96OR22464 with Lockheed Martin Energy System Research Cooperation.

Abstract

A new apparatus has been designed and built. It includes a modified commercial double-pass cylindrical mirror analyzer (CMA) and several time of flight (TOF) spectrometers. Energy and angular resolved electron emission spectra can be obtained with this apparatus and electron-electron and electron-ion coincidence measurements can be made. The apparatus has been used in three different experiments.

Electron-emission yields from the doubly excited states $2s^2(^1S)$, $2p^2(^1D)$, $2s2p(^1P)$ and $2p^2(^1S)$ of helium produced by 1.67MeV/u carbon ions with charge states $q=4,5,6$ were measured with an energy resolution better than accomplished before. A CMA was used to observe simultaneously electrons ejected between 47.7° and 132.3° with respect to the ion-beam direction. Comparison of the data with several calculations was made to understand the mechanisms responsible for the production of these states. The relative electron emission yields for the $2p^2(^1D)$ and $2s2p(^1P)$ states were studied as a function of the projectile charge state.

Electron-electron coincidence measurements from Auger-electron cascade decay processes after photon atom interactions were also performed. When a 3203.5eV photon interacts with an argon atom, producing a 1s hole, the atom typically decays by KLL Auger electron emission so that two holes in the L -shell are created; the atoms can further decay by LMM Auger electron emission. The result is a cascade-like decay process. Angular resolved coincidence measurements were made between the emitted KLL and the LMM Auger electrons. These coincidence measurements allow the study of the angular correlation between the emission of different first step KLL Auger electrons with the second step LMM Auger electron group.

When cascade decays take place, the photoion (in this case the argon ion) final charge state will depend on the decay path. The measured charge state fractions, however, depend on several experimental parameters. A detailed study of the variation of the measured charge state fractions with the experimental parameters was

made. Coincidence measurements between KLL Auger electrons and recoil photoions were made at different photon energies. Branching ratios for several decay processes were obtained.

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1 Introduction

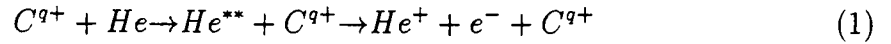
Three distinct experiments were carried out for this thesis. We have measured the electron yield from the most prominent doubly excited states of He produced by fast carbon ion impact, as a function of the electron energy and emission angle. We have also measured the angular correlation between the Ar *KLL* and *LMM* Auger electrons coming from cascade decay processes after *K*-shell photoexcitation. In addition, we have studied how the experimental parameters can affect measurements of the Ar photoion charge-state distribution arising from such cascades. After improving control over the experimental parameters, we made non-coincidence measurements of the Ar photoion charge-state distribution as a function of the photon energy, and measurements of the Ar photoion charge-state distribution observed in coincidence with the Ar *KLL* Auger electrons. We will give an introduction for all three experiments.

1.1 Doubly Excited States of Helium Produced by Fast-Ion Impact

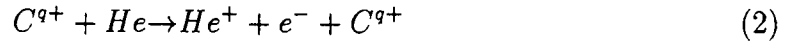
The electron-emission yields from the doubly excited states $2s^2(^1S)$, $2p^2(^1D)$, $2s2p(^1P)$ and $2p^2(^1S)$ of helium produced by 1.67MeV/u carbon ions with charge states $q = 4, 5, 6$ are measured with an energy resolution better than previously accomplished. A cylindrical mirror analyzer (CMA) is used to observe simultaneously the electrons ejected between 47.7° and 132.3° with respect to the ion beam direction. Comparison of the data with several calculations is made to understand the mechanisms responsible for the production of these states. The relative electron-emission yields for the production of the $2p^2(^1D)$ and $2s2p(^1P)$ states are studied as a function of the projectile charge state.

During recent years there has been increased interest in studying electron-correlation effects in atomic collisions. Studies of double ionization [1, 2] and transfer

ionization [3, 4] of He by fast bare projectiles demonstrated the importance of the electron-electron interaction. These effects have been studied in the doubly excited states of helium produced by fast projectile ions by Pedersen *et al.* [5] and Giese *et al.* [6],[7]. Using 20MeV carbon ion impact we measured the electron emission from doubly excited states in helium into a wide angular range at a resolution better than achieved before, allowing us to better study these effects. The possible two-electron phenomena in the collision of bare projectiles with He include double excitation, excitation ionization, double ionization, transfer excitation, transfer ionization, and double transfer. The two processes that contribute to our electron-emission spectra that follow from the interaction between the carbon ions and the helium target atoms are double excitation followed by autoionization



that appears in the energy spectra as a resonance; and direct ionization



that appears as a continuum background. The autoionized electrons can interfere with directly ionized electrons of the same energy. As a result, the energy profiles and angular distributions of emission lines due to autoionizing states can be asymmetric. A schematic drawing of the process studied is shown in Fig. 1, and a diagram of energy levels of He is shown in Appendix C.

One of the simplest models used to describe the double-excitation process in detail is the independent-electron model (IEM). The IEM assumes that each target electron interacts independently with the projectile [8], [9]. For two-electron excitation this is known as a second-order process [10] (TS-2). It does not include any correlation between the electrons. The total cross section for this process is expected to vary as q^4 , where q is the projectile charge state. There are also first-order mechanisms [6][10] (TS-1) which are used to describe the double-excitation process. These are roughly

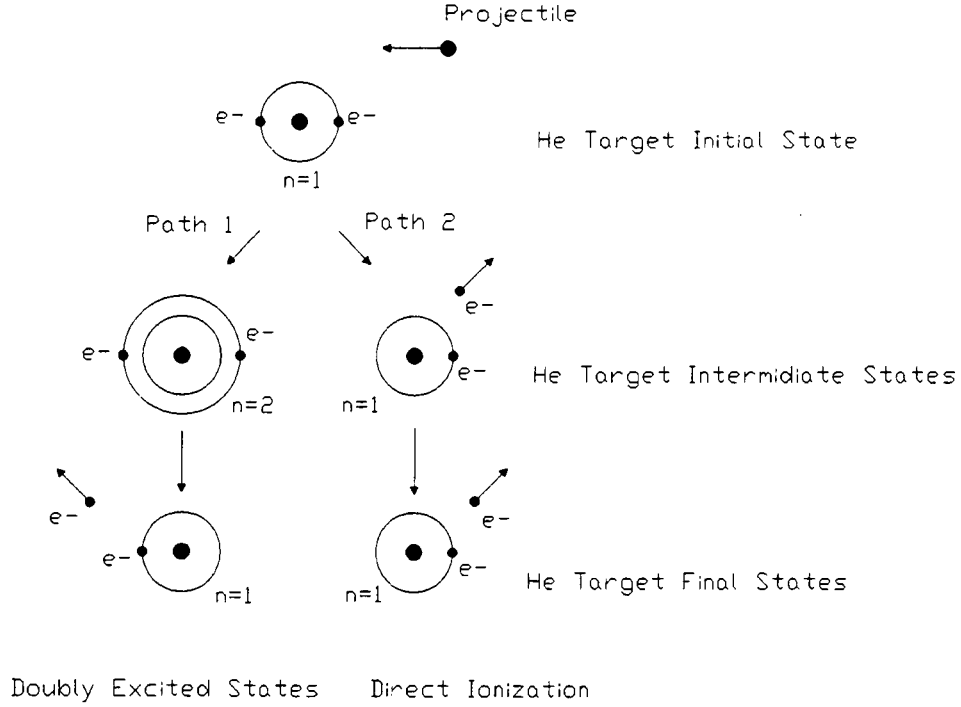


Figure 1: Schematic drawing of the two processes that contribute to the electron-emission yield.

divided into two types [6], static, and dynamic correlation of the target electrons. The static mechanism assumes that the projectile interacts with only one target electron. The second electron is excited due to a final state rearrangement of the target produced by a static correlation of the electrons, i.e. a “shake up” process. The dynamic mechanism also assumes that the projectile interacts with only one electron, but now this electron can, in an oversimplified description, collide with the remaining electron resulting in a two-step process. The total cross section for these two first-order mechanisms is expected to vary as q^2 . A quantum interference between the first- and second-order mechanisms will give a term in the total cross section that is proportional to q^3 .

More recent calculations done by Morishita *et al.* [11] allow us to directly com-

pare the measured electron-emission yields with the calculated ones. Calculations done by Martín *et al.* [12] give us information about the peak shape of the autoionized electron interference with the directly ionized electrons. In the work done by Morishita *et al.* [11] the authors assume that due to the high impact energy the projectile exerts a pulse-like force on the He system, so that by exploiting the accurate solution of the three-body subsystem, He, and by treating the projectile as a classically moving source of the external field, they can calculate closely the profile of each resonance against the continuum background. The principle of the method for the calculation done by Morishita *et al.* [11] may be summarized as follows:

- 1) The projectile motion is represented by a classical straight-line trajectory with a well-defined value of the impact parameter, and is thus described as a moving source of the external field;
- 2) The He target state wave functions are calculated by the hyperspherical close-coupling (HSCC) method. Thus, the autoionization mechanism is properly contained in the continuum-state wave functions;
- 3) Wave-packet-type energy-discretization of the target continuum states is effected on a flexibly chosen energy mesh so as to evaluate the resonance profile of each resonance closely; and
- 4) The time-dependent Schrödinger equation is thus obtained and solved by the close-coupling (CC) method.

The calculations performed by Martín *et al.* [12] involved a single-center close-coupling calculation, correlated two-electron states, and a coarse-grain discretization of the continuum, and was done in the impact-parameter framework. The basis used in their calculation include the $2s^2(^1S)$, $2s2p(^1P)$ and $2p^2(^1D)$ doubly excited states of He, the ground state of He, all bound excited states of He ($1s2s$, $1s2p$, $1s3s$, ...) for $L = 0, 1, 2$; all discretized continuum states with $E \leq 35au$ ($10au$ for $L = 2$) for $L = 0, 1, 2$; and the discretized continuum states which are degenerate with the ^1D resonance, for $L = 3, 4$. All the M values associated with a given L were kept in the

calculation.

1.1.1 Autoionization Theory

The experimental observation of an autoionized electron measures a physical state which can be defined by the electron energy and the emission angle, $|E_e, \theta\rangle$. This physical state was described by Fano [13] as a mixture of a discrete doubly excited state $|R\rangle$ and a continuum of single ionization states $|E_e^0, \theta\rangle$. The states $|R\rangle$ and $|E_e^0, \theta\rangle$ constitute a basis set that is required to be complete and orthogonal and that the transformation from the physical state to this basis set be unitary. The physical state can be formally represented as

$$|E_e, \theta\rangle = |R\rangle\langle R|E_e, \theta\rangle + \int dE_e^0 |E_e^0, \theta\rangle\langle E_e^0, \theta|E_e, \theta\rangle \quad (3)$$

where the parameters $\langle R|E_e, \theta\rangle$ and $\langle E_e^0, \theta|E_e, \theta\rangle$ are the expansion coefficients of the physical state on this basis set.

The measured electron-emission yield can be expressed as

$$\frac{d^2\sigma}{dE_e d\Omega_e} = |\langle E_e, \theta|T|i\rangle|^2 \quad (4)$$

where $|i\rangle$ is the initial ground state of the He, and T represents the perturbation caused by the projectile. The yield in the terms of the common basis set is

$$\frac{d^2\sigma}{dE_e d\Omega_e} = \left| \langle E_e, \theta|R\rangle\langle R|T|i\rangle + \int dE_e^0 \langle E_e, \theta|E_e^0\rangle\langle E_e^0, \theta|T|i\rangle \right|^2 \quad (5)$$

The two terms on the right hand side of this equation can interfere. The strength of the interference depends on the relative magnitudes of these two terms and on the relative phase between them. There is no *a priori* reason for the interference to be constructive or destructive (positive or negative). It is this interference between the amplitude for the double excitation $\langle R|T|i\rangle$ and the amplitude for the single

ionization $\langle E_e^0, \theta | T | i \rangle$ that makes the extraction of the double-excitation cross section from the data difficult.

A common parametrization of the doubly differential cross section for emission of electrons in the region of an isolated autoionizing resonance is called the Shore parametrization [14],

$$\frac{d^2\sigma}{dE_e d\Omega_e} = f(E_e, \theta) + \frac{a(\theta)\epsilon + b(\theta)}{\epsilon^2 + 1} \quad (6)$$

where

$$\epsilon = \frac{2(E_e - E_r)}{\Gamma_r} \quad (7)$$

E_r is the energy of the resonance and Γ_r is the natural width of the resonance. The incoherent contribution to the differential cross section from the direct ionization is given by $f(E_e, \theta)$, the asymmetry of the profile is described by $a(\theta)$ and the contribution of the discrete state to the cross section is given by $b(\theta)$. The latter can be illustrated by integrating over the electron energy to obtain the singly differential cross section

$$\frac{d\sigma}{d\Omega_e} = F(\theta) + \frac{1}{2}\pi\Gamma_r b(\theta) \quad (8)$$

where $F(\theta)$ now represents the incoherent contribution from the direct ionization.

The differential emission yield, $d\sigma/d\Omega_e$, is approximately equal to the differential single-ionization yield for energies far away from the resonance energy, that is for $\epsilon \gg 1$ it is true that

$$\frac{d^2\sigma}{dE_e d\Omega_e} \simeq f(E_e, \theta) \equiv \frac{d^2\sigma_{SI}}{dE_e d\Omega_e} \quad (9)$$

Further, it can be reasonably assumed that the single-ionization yield would, in the absence of a resonance, vary smoothly with energy. The parameters $a(\theta)$ and $b(\theta)$ can in principle be positive or negative. However, depending on the relative magnitudes of $a(\theta)$ and $b(\theta)$, the resonance can have any shape from a pure dip (or window) in

the background to a pure Lorentzian peak on top of the background. The quantity $b(\theta)$ measures the intensity of the resonance and $a(\theta)$ measures its shape.

The peak profile can also be represented by the Fano parametrization [13], which is written as

$$\frac{d^2\sigma}{dE_e d\Omega_e} = f_{back}(E_e, \theta) + f_{res}(\theta) \frac{[\tilde{q}(\theta) + \epsilon]^2}{\epsilon^2 + 1} \quad (10)$$

where $f_{back}(E_e, \theta)$ represents the transitions to states of the continuum which do not interact with the discrete state, $f_{res}(\theta)$ represents the transitions to states which do interact with the discrete states, and ϵ is defined as in equation (7). The dimensionless parameter $\tilde{q}$ is a measure of the asymmetry of the profile but is independent of the magnitude of the cross section. Note that $\tilde{q}$ is not the same as the Fano shape parameter [13]. The latter was derived to describe the shape of the resonance in the projectile scattering angle and final energy. The relationship of $\tilde{q}$ to the Shore parameters is given by

$$\tilde{q} = \frac{b(\theta) + [b(\theta)^2 + a(\theta)^2]^{\frac{1}{2}}}{a(\theta)} \quad (11)$$

When the absolute value of $\tilde{q}$ is very large, the peak is nearly Lorentzian. When $\tilde{q}$ is approximately zero, the peak becomes a symmetric dip or window in the continuum. The profile is asymmetric for intermediate values of $\tilde{q}$, with negative values corresponding to profiles with a dip on the high-energy side of the resonance, and positive values to a dip on the low-energy side. The Shore parametrization and the Fano parametrization are equivalent. We have chosen the Shore parametrization to fit our measured spectra, because $b(\theta)$ is proportional to the contribution of the double-excitation resonance to the differential ionization cross section. Thus $b(\theta)$ can be used to compare the excitation by different projectiles.

The relationship between $b(\theta)$ and the double-excitation cross section can be found by integrating equation (8) over the electron emission angle to obtain the total emission cross section

$$\sigma = \sigma_{SI} + \frac{\pi\Gamma_r}{2} \int b(\theta) d\Omega_e = \sigma_{SI} + \frac{\pi\Gamma_r}{2} B. \quad (12)$$

The total cross section in terms of the preferred basis set is

$$\begin{aligned} \sigma &= \iint d\Omega_e dE_e |\langle E_e, \theta | T | i \rangle|^2 \\ &= \iint d\Omega_e dE_e \left| \langle E_e, \theta | R \rangle \langle R | T | i \rangle + \int dE_e^0 \langle E_e, \theta | E_e^0, \theta \rangle \langle E_e^0, \theta | T | i \rangle \right|^2 \end{aligned} \quad (13)$$

Using the orthonormality and completeness of the basis set, this expression can be simplified to

$$\iint d\Omega_e dE_e |\langle E_e, \theta | T | i \rangle|^2 = \iint d\Omega_e dE_e |\langle E_e^0, \theta | T | i \rangle|^2 + |\langle R | T | i \rangle|^2 \quad (14)$$

The last term in this equation $|\langle R | T | i \rangle|^2$ is the double-excitation cross section. From equation (8) it is possible to see that $b(\theta)$ is proportional to the single differential emission yield minus the extrapolated single-ionization background. Rearranging equation (14) gives

$$\int d\Omega_e \left[dE_e \left(|\langle E_e, \theta | T | i \rangle|^2 - |\langle E_e^0, \theta | T | i \rangle|^2 \right) \right] = \iint d\Omega_e dE_e |\langle R | T | i \rangle|^2 \equiv \sigma_{DE} \quad (15)$$

where the term in large parentheses is proportional to $b(\theta)$. Equation (15) is important because it gives the relationship between the measured electron-emission yield and the calculated single-ionization and double-excitation cross sections. Further, equation (15) shows that $B = \int b(\theta) d\theta$ is directly proportional to the double-excitation cross section. Note that while $b(\theta)$ can be positive or negative, the integral of $b(\theta)$ over the electron energy and emission angle represents a cross section and must be positive. If $b(\theta)$ is negative at some angles, it must therefore become positive at other angles in order for B to be positive.

The above analysis shows that it is, in principle, possible to remove these interference effects and obtain the double-excitation cross section by integrating the

difference between the measured electron-emission yield and the extrapolated single-ionization yield over both the electron energy and the electron-emission angle. However, due to the small separation of the $2p^2(^1D)$ and $2s2p(^1P)$ states, it is not possible to obtain individual cross sections for these two overlapping states by this process. Instead, a more detailed analysis of the line shape for each state is necessary, using, for example, the Shore parametrization as will be shown in Section 3.1.

1.2 Cascade-Decay Processes Produced in Photon-Atom Interactions

When a vacancy is produced in an inner shell of a many-electron atom (e.g., by photoionization or photoexcitation), a rearrangement of the system usually proceeds through a stepwise series of radiative (i.e., photon emission) and non-radiative (i.e., Auger electron emission) decay processes known as a vacancy cascade decay [15, 16], as is schematically shown in Fig. 2 for *Ar*. This process results in the emission of X-ray (fluorescence) photons and Auger electrons and often leads to a highly charged residual ion. Due to the multitude of decay pathways two-parameter coincidence measurements such as Auger electron-ion [21]; threshold photoelectron-residual ion [20]; Auger electron-Auger electron or Auger electron-residual ion [17, 18, 19] measurements, can help in the analysis of the vacancy cascade decay processes. The advantage of using coincidence techniques is that a subset of transitions are selected resulting in simplified electron energy spectra or ion charge state distributions. Several theoretical calculations done by Kabachnick *et al.* [16, 22], and El-Shemi *et al.* [23], analyze the atomic reorganization cascades that follow an inner-shell vacancy.

In the present work we made Auger electron-Auger electron coincidence measurements of argon atoms photoexcited with a linearly polarized photon beam with an energy equal to the *Ar* $1s \rightarrow 4p$ resonance (3203.54 eV). Some of the possible cascade processes that can occur from this interaction are shown in Table 1. In Table 1 the

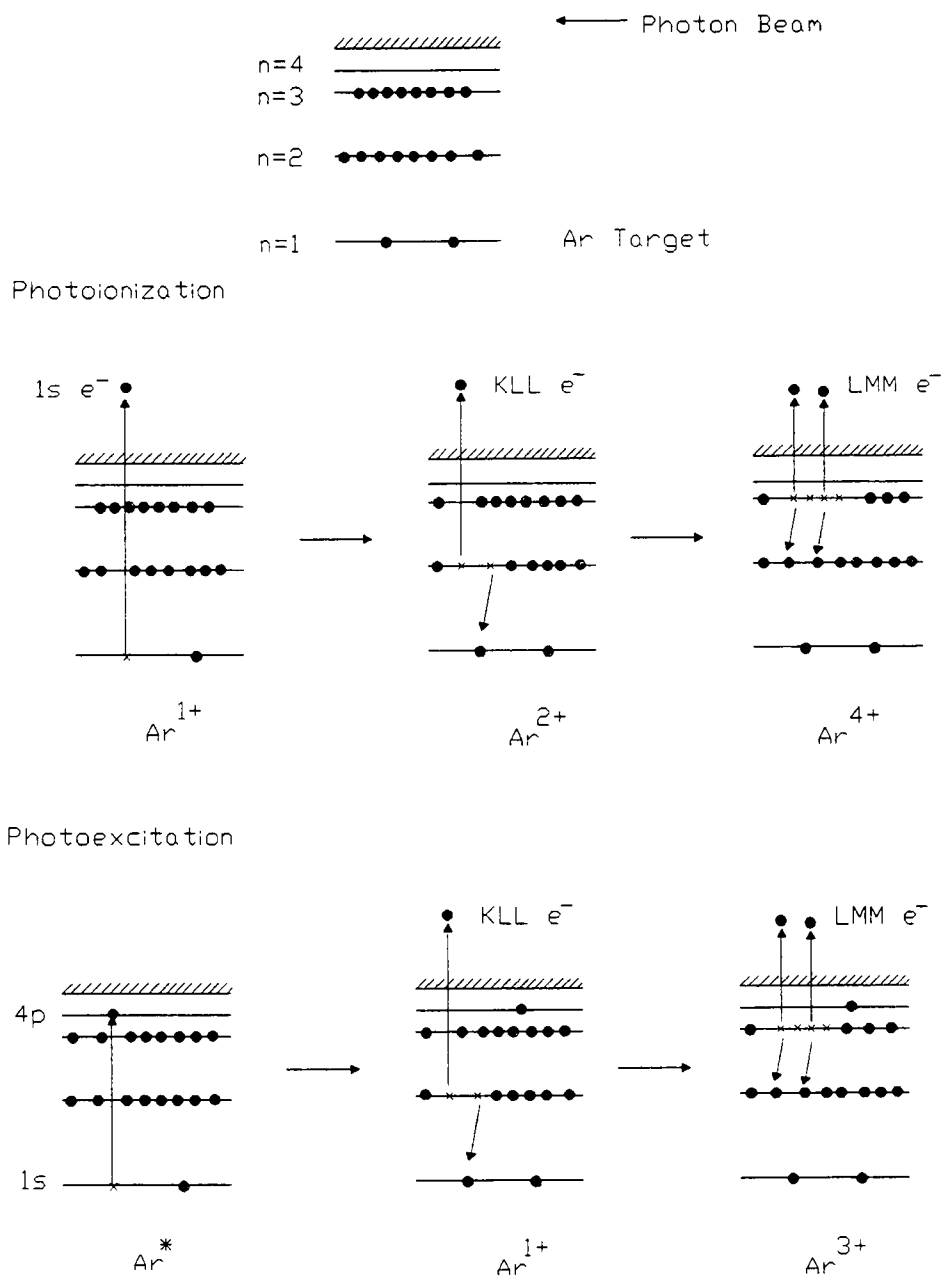


Figure 2: Illustration of *Ar* *K*-shell photoionization and photoexcitation followed by $K \rightarrow L \rightarrow M$ vacancy cascades. Electrons are indicated by circles and holes by crosses.

Table 1: *Ar* Auger-decay processes of *K*-holes produced by photoionization or photoexcitation (with the *4p* electron as a spectator).

Auger decay	Subsequent processes	Subsequent processes	Produced ions photoionization	Produced ions photoexcitation
$KL_{23}M$	$L_{23}MM$		Ar^{3+}	Ar^{2+}
$KL_{23}L_{23}$	$2(L_{23}MM)$		Ar^{4+}	Ar^{3+}
KL_1M	L_1MM		Ar^{3+}	Ar^{2+}
	$L_1L_{23}M$	$L_{23}MM$	Ar^{4+}	Ar^{3+}
KL_1L_{23}	$L_{23}MM + L_1MM$		Ar^{4+}	Ar^{3+}
	$L_{23}MM + L_1L_{23}M$	$L_{23}MM$	Ar^{5+}	Ar^{4+}
KL_1L_1	$2(L_1MM)$		Ar^{4+}	Ar^{3+}
	$L_1MM + L_1L_{23}M$	$L_{23}MM$	Ar^{5+}	Ar^{4+}
	$2(L_1L_{23}M)$	$2(L_{23}MM)$	Ar^{6+}	Ar^{5+}

residual ion charge states in the case where the photon energy is tuned to the *Ar* $1s \rightarrow 4p$ resonance, or alternately above the *Ar* *K*-threshold at 3206.26eV, are also shown. We concentrated on coincidence measurements between the $KL_{23}L_{23}$ Auger electrons as the first step in the Auger-cascade decay, and the subsequently emitted *LMM* Auger electrons as the second. Due to our capability of measuring the angular distribution of $KL_{23}L_{23}$ Auger electrons we were able to study the possible angular correlation between two Auger electrons from the cascade decay process. We also looked at KL_1L_{23} and $KL_{23}M$ Auger electrons as the first step Auger decay, in coincidence with subsequent emitted *LMM* Auger electrons. As the intensity of these two Auger transitions is much smaller than for the $KL_{23}L_{23}$ processes we were not able to obtain detailed angular distributions for them.

From the angle-resolved coincidence measurements it is possible to determine, for two Auger electrons emitted in a multi-step decay process from an inner shell vacancy,

the particular angular momenta of the atomic state involved in the transition. These measurements can then be used to determine the angular momenta and parity of the particular lines in the complex Auger spectrum.

1.2.1 Theory of Angular Correlations in Cascades

Let us consider the simplest case of two successive decays (either an Auger electron or x-ray photon emission) of an excited atom. We assume that the atomic transitions occur in steps and the atomic states involved at each stage have a well defined total angular momentum and parity. The initial state with total angular momentum J_0 may be oriented and/or aligned. It is convenient to describe the anisotropy of the state by the statistical tensor $\rho_{k_0 q_0}(J_0, J_0)$ [24]. In general, all the components of the statistical tensor with $k_0 \leq 2J_0$, $-k_0 \leq q_0 \leq k_0$ can be non-zero. If the initial state is produced by photoionization or photoexcitation of an inner-shell electron with unpolarized or linearly polarized light, then the vacancy state is in general aligned and the degree of alignment is described by a second-rank statistical tensor ($k_0 = 2$). If we also assume that the detectors are insensitive to the polarization (spin) state of the radiation, then the general expression for the angular-correlation function of two successive decays of the polarized initial state can be expressed in the following form:

$$W(\mathbf{n}_1, \mathbf{n}_2) = 4\pi \hat{J}_0 \sum_{k_0 q_0 k_1 k_2} G_{k_1 k_2 k_0} \rho_{k_0 q_0}(J_0, J_0) [Y_{k_1}(\mathbf{n}_1) \otimes Y_{k_2}(\mathbf{n}_2)]_{k_0 q_0} \quad (16)$$

Here $\mathbf{n}_1$ and $\mathbf{n}_2$ are the unit vectors which determine the direction of the first and second radiation, and $\hat{J} = \sqrt{2J+1}$. The tensorial product of the spherical functions is defined as

$$[Y_{k_1}(\mathbf{n}_1) \otimes Y_{k_2}(\mathbf{n}_2)]_{k_0 q_0} = \sum_{q_1 q_2} (k_1 q_1, k_2 q_2 | k_0 q_0) Y_{k_1 q_1}(\theta_1, \phi_1) Y_{k_2 q_2}(\theta_2, \phi_2) \quad (17)$$

where $(k_1 q_1, k_2 q_2 | k_0 q_0)$ is a Clebsch-Gordan coefficient.

The generalized angular-correlation coefficients $G_{k_1 k_2 k_0}$ contain the angular momentum coupling coefficients and amplitudes of the decay process. As we assume that the decay process occurs in steps, each of the terms $G_{k_1 k_2 k_0}$ may be split into two factors, each of which depends entirely on the corresponding transition, then

$$G_{k_1 k_2 k_0} = B_{k_1 k_2 k_0}(J_0, J_1) A_{k_2}(J_1, J_2). \quad (18)$$

Here J_1 and J_2 are the total angular momenta of the intermediate and the final atomic states, respectively. These factors are given by

$$B_{k_1 k_2 k_0}(J_0, J_1) = (-1)^{k_1+k_2-k_0} \hat{J}_0 \hat{J}_1 \sum_{L_1 L_1'} c_{k_1 0}^*(L_1, L_1') \times \begin{Bmatrix} J_1 & L_1 & J_0 \\ J_1 & L_1' & J_0 \\ k_2 & k_1 & k_0 \end{Bmatrix} \langle J_1, L_1 || \hat{O} || J_0 \rangle \langle J_1, L_1' || \hat{O} || J_0 \rangle^*, \quad (19)$$

$$A_{k_2}(J_1, J_2) = (-1)^{J_1+J_2+k_2} \hat{J}_1 \sum_{L_2 L_2'} c_{k_2 0}^*(L_2, L_2') (-1)^{L_2'} \times \begin{Bmatrix} J_1 & L_2 & J_2 \\ L_2' & J_1 & k_2 \end{Bmatrix} \langle J_2, L_2 || \hat{O} || J_1 \rangle \langle J_2, L_2' || \hat{O} || J_1 \rangle^*. \quad (20)$$

Here we used the standard notation for the Wigner $6j$ and $9j$ coefficients; $\langle J_b, L || \hat{O} || J_a \rangle$ denotes the decay amplitude which describes a transition from the state J_a to the state J_b by emission of the radiation with the total angular momentum L . The $c_{k_0}(L, L')$ are radiation parameters. For an Auger electron the radiation parameters are

$$c_{k_0}(L, L') = (-1)^{L'+l+\frac{1}{2}} (4\pi)^{-1} \hat{l} \hat{l}' \hat{L} \hat{L}' (l 0, l' 0 | k 0) \begin{Bmatrix} l & L & \frac{1}{2} \\ L' & l' & k \end{Bmatrix} \quad (21)$$

where L and l are the total and orbital angular momenta of the Auger electron respectively (in this case the summation in (19) and (20) include also the sums over l and l' .) For the dipole fluorescence (the detector is not sensitive to polarization) the radiation parameters are

$$c_{k_0}(1, 1) = \frac{3}{8\pi}(11, 1 - 1|k_0). \quad (22)$$

Since the atomic states involved are assumed to have well-defined parities and parity is conserved in the decay processes, it can be shown that only even values of k_1 and k_2 can occur in (16). In addition the summation in (16) is limited by the triangle rule $|k_0 - k_1| \leq k_2 \leq k_0 + k_1$.

Expression (16) shows that each of the terms under summation consists of two parts. The first factor, the angular-correlation coefficient, is determined by the dynamics of the particular transition, while the second factor arise purely from kinematics; it is entirely determined by the geometry of the experiment and the initial conditions. The second factor is common for any emitted electron or photon. The main goal of a correlation experiment is to obtain the angular-correlation coefficients and as a final result to determine the amplitudes of the process. However, the analysis of the kinematics factor in (16) is necessary in order to match the experimental conditions.

By choosing a particular geometry for the experiment, the angular-correlation function can be simplified, making the analysis of the experimental data easier. As the creation of the initial vacancy in our experiment is produced by synchrotron radiation (SR), we consider here that the photoexcitation or photoionization of the core electron is produced by linearly polarized light. We choose the Z-axis along the direction of the linear polarization. The vacancy or core-excited initial state is in general aligned along this direction. If one of the cascade products is detected along the polarization direction, the experimental conditions for the second radiation turns out to be axially symmetric with respect to this direction. Therefore the

angular correlation will depend only on one polar angle θ . For this case the angular correlation can be obtained from equation (16), where

$$W(\theta) = \sum_{k=\text{even}} a_k P_k(\cos \theta), \quad (23)$$

and

$$a_k = \hat{J}_0 \sum_{k_0 k_1} \hat{k}_0 \hat{k}_1 (k_1 0, k_0 | k_0 0) \rho_{k_0 0}(J_0, J_0) G_{k_1 k k_0}. \quad (24)$$

If the two measured decay products are Auger or autoionization electrons ($e - e$ correlation) the complexity of the correlation function depends on their orbital angular momenta. In particular, if the initial vacancy is unpolarized, as in the ionization of an inner s -shell, then $k_0 = q_0 = 0$, and from equation (17) it follows that $k_1 = k_2$ and expression (16) takes the form:

$$W(\theta) = \sum_{k=\text{even}} \bar{A}_k(J_0, J_1) A_k(J_1, J_2) P_k(\cos \theta) \quad (25)$$

where $P_k(\cos \theta)$ are the Legendre polynomials; the factors $A_k(J_1, J_2)$ are defined in (20). For $\bar{A}_k(J_0, J_1)$ one obtains

$$\begin{aligned} \bar{A}_k(J_0, J_1) = & (-1)^{J_1+J_0} \hat{J}_1 \sum_{L_1 L'_1} c_{k_0^*}(L_1, L'_1) \\ & \times (-1)^{L_1} \left\{ \begin{matrix} J_1 & L_1 & J_0 \\ L'_1 & J_1 & k \end{matrix} \right\} \langle J_1, L_1 || \hat{O} || J_0 \rangle \langle J_1, L'_1 || \hat{O} || J_0 \rangle^*. \end{aligned} \quad (26)$$

Now the angular-correlation function depends only on one angle θ , the relative angle between the two electrons. If the initial state is unpolarized, then the first electron is emitted isotropically. However, if we fix the direction of emission and if the electron takes a definite angular momentum from the atom, then the intermediate atomic state will be aligned along the direction of the electron radiation. The second electron can emitted be anisotropic, but its angular distribution will be axially symmetrical about the direction of the first electron. For even k , the distribution is also symmetric

with respect to reflection through the $\theta = 90^\circ$ plane. The complexity of the angular correlation is determined by the condition $k \leq \min(2L_1, 2L_2, 2J_1)$ where L_1 and L_2 are the total angular momenta of the two electrons. From this result we can see that if either electron or the intermediate state has angular momentum equal to 0 or $\frac{1}{2}$, the angular correlation disappears and $W = \text{const.}$

Thus, from the angle-resolved coincidence measurements, it is possible to determine, for two Auger electrons emitted in a multi-step cascade decay process from an inner-shell vacancy, the particular angular momenta of the atomic state involved in the transition. These measurements can then be used to determine the angular momenta and parity of the particular lines in the complex Auger spectrum, for the cases where the spectral resolution is sufficient to permit this determination.

1.2.2 Basic Relations and Background

Synchrotron radiation (SR) is a very important tool for electron spectrometry studies, due to the wide and continuously variable photon-energy range, the high photon flux and the narrow energy bandpass achievable with SR. In electron spectrometry, the kinetic energies of electrons ejected from an atom or molecule by photon, charged particles, an Auger process, or some other excitation mode are usually measured [25]. Since photoionization played a central role in electron spectrometry using synchrotron radiation (ESSR), some basic relations for the photoeffect are presented here.

The Einstein photoeffect relation forms the basis of photoelectron spectrometry by interrelating the photon energy, $h\nu$, with the kinetic energy of the photoelectron, $E_{e,kin}$, ejected from the atomic or molecular level with binding energy E_B

$$h\nu = E_{e,kin} + E_B \quad (27)$$

A more detailed description of the interaction of a photon γ (with energy $h\nu$, angular momentum j , and parity π) with an atom A (with total ground-state energy E_0 , total

angular momentum J_0 and parity π_0) is provided by an expression of the form

$$\begin{aligned} & \gamma(h\nu, j_\gamma = 1, \pi_\gamma = -1) + A(E_0, J_0, \pi_0) \\ & \rightarrow A^+(E_f, J_f, \pi_f) + e^-(E_{e,kin}, l s j, \pi_e = (-1)^l), \end{aligned} \quad (28)$$

where the ionized atom is characterized by the final-state parameters E_f , J_f , π_f and the outgoing photoelectron by $E_{e,kin}$, $l s j$, and π_e . In this reaction we consider that a single electron is emitted from an atom in an electric dipole interaction. Since energy, angular momentum and parity are conserved, then

$$\begin{aligned} E_{e,kin} &= h\nu - (E_0 - E_f) \\ \mathbf{J}_0 + \mathbf{j}_\gamma &= \mathbf{J}_f + \mathbf{s} + \mathbf{l} \\ \pi_0 \cdot \pi_\gamma = \pi_f \cdot \pi_e &= -\pi_0 = (-1)^l \pi_f \end{aligned} \quad (29)$$

The photoionization cross section can be written as [26, 25]

$$\sigma_{if} = \frac{4\pi^2\alpha^2}{k_\nu} \left| \langle \Psi_i^* | \sum_\mu \exp(-i\mathbf{K}_\nu \cdot \mathbf{r}_\mu) \epsilon \nabla_\mu | \Psi_f \rangle \right|^2 \quad (30)$$

where α is the fine-structure constant, $\mathbf{k}_\nu$ and ϵ are the momentum and polarization vectors of the incident photon, $\mathbf{r}_\mu$ and ∇_μ are the position and gradient operators of the μ th electron, and Ψ_i and Ψ_f represent the initial (ground) state and final state with their wave function properly normalized. As the continuum electron, which is part of the final state Ψ_f , can be specified by its direction, equation (30) also represents the formula for the differential cross section $d\sigma_{if}/d\Omega$. Equation (30) shows that the cross section depends on the multipolarity of the photon field through $\exp(-i\mathbf{K}_\nu \cdot \mathbf{r}_\mu)$, on the polarization, and on the initial- and final-state wave functions of the system through Ψ_i and Ψ_f . Theory has concentrated on calculating the best Ψ_i and Ψ_f with fewest approximations in the solution of the Schrödinger or

Dirac equation. The frequently used dipole approximation (neglecting retardation) assumes that $\exp(-i\mathbf{K}_\nu \cdot \mathbf{r}_\mu) = 1$ because $|\mathbf{K}_\nu \cdot \mathbf{r}_\mu| \ll 1$.

Then based on equation (30) [25] the photoelectron angular distribution can be characterized by coefficients B_n in a Legendre expansion [27, 28] as

$$\frac{d\sigma_{if}}{d\Omega} = \frac{\sigma_{if}}{4\pi} \sum_{n=0}^{\infty} B_n P_n(\cos \theta) \quad (31)$$

where $P_n(\cos \theta)$ are the Legendre polynomials. In the nonrelativistic dipole approximation all the coefficients B_n are zero except for B_2 ($B_0 \equiv 1$ always). Then B_2 is the single parameter that determines the angular distribution. Assuming polarized radiation and a randomly oriented target,

$$\frac{d\sigma_{if}}{d\Omega} = \frac{\sigma_{if}}{4\pi} \left[1 + \frac{\beta_{if}}{2} (3 \cos^2 \theta - 1) \right], \quad (32)$$

where θ is the angle between the photoelectron direction and the polarization vector $\mathbf{E}$, $P_2(\cos \theta) = (3 \cos^2 \theta - 1)/2$ and $B_2 = \beta_{if}$, where β_{if} is the asymmetry parameter. The photoelectron angular distribution for several values of β_{if} are shown in Fig. 3. In Fig. 3, θ_m is known as the magic angle, for which the value of $\frac{d\sigma_{if}}{d\Omega}$ is the same for all β_{if} .

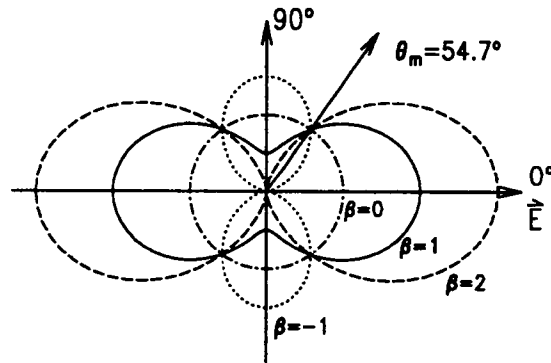


Figure 3: Photoelectron angular distribution for different values of the asymmetry parameter β .

1.3 Photoion Charge-State Distributions and Electron-Ion Coincidences

We measured the photoion charge state distribution as a function of the photon energy in the vicinity of the *Ar* *K*-threshold. Measurements of photoions in coincidence with different *KLL* Auger electrons were also conducted, for photon energies at the *Ar* $1s \rightarrow 4p$ resonance and $10eV$ above resonance. The multiply charged ions are produced mainly by successive Auger decay processes after the initial vacancy is formed, as it is schematically shown in Fig. 2 in Section 1.2. Another possibility is shake-off of electrons accompanying the vacancy creation, with or without recapture to an excited state due to post collision interactions (PCI) with the slowly moving photoelectron formed when the excitation is near threshold. Another possibility is the elimination of two or more electrons due to direct processes, which are possible because the electrons in the atom are not independent, but correlated [29]. Decays by photon emission are also possible, but such decays only produce singly and doubly charged ions, as will be discussed later [30].

Several measurements of the photoion charge state distribution near the argon *K*-edge are available in the literature [31, 17, 32], but recent measurements of the double to single ionization ratio of helium showed how critical certain experimental parameters are to obtain the correct values for the ratio [33]. Any experiment where the detection of ions is performed with a microchannel plate detector (MCP) requires careful control of these critical experimental parameters. We present here a detailed study of these photoion charge state distributions as a function of the critical experimental parameters that affect the experimental results. After the optimum experimental parameters were determined we measured the photoion charge state fractions as a function of the photon energy so we were able to obtain the branching ratios for different steps in the cascade decay process. Coincidence measurements of the photoions with the argon $KL_{23}L_{23}$ and KL_1L_{23} Auger electrons were done at

two photon energies, 3203.5eV and 3213.5eV . We were able to obtain the branching ratios for the second step processes after selecting a given first step decay path.

Some of the possible non-radiative decay paths are shown in Table 1 in Section 1.2, where the final ion charge states are indicated for the photoionization and photoexcitation cases, for photon energies near and above the K -shell threshold. From this table one can also see that by performing a coincidence measurement the distribution of the charge states will be dramatically simplified. The most probable decay paths for a first step $KL_{23}L_{23}$ and KL_1L_{23} Auger electrons, for photon energies near and above the K -shell threshold, are the following:

$$\begin{aligned}
KL_{23}L_{23} &\longrightarrow L_{23}MM + h\nu[M] \longrightarrow Ar^{3+} \\
KL_{23}L_{23} &\longrightarrow 2(L_{23}MM)n'_p \longrightarrow Ar^{3+} \\
KL_{23}L_{23} &\longrightarrow 2(L_{23}MM) \longrightarrow Ar^{4+} \\
KL_{23}L_{23} &\longrightarrow (2(L_{23}MM) + MM)n'_p \longrightarrow Ar^{4+} \\
KL_{23}L_{23} &\longrightarrow 2(L_{23}M_1M_1) + M_1M_{23} \longrightarrow Ar^{5+}
\end{aligned} \tag{33}$$

$$\begin{aligned}
KL_1L_{23} &\longrightarrow L_{23}MM + L_1MM \longrightarrow Ar^{4+} \\
KL_1L_{23} &\longrightarrow L_{23}MM + L_1MM \rightarrow (M_1M_{23})n'_p \longrightarrow Ar^{4+} \\
KL_1L_{23} &\longrightarrow L_{23}MM + L_1L_{23}M \rightarrow L_{23}MM \longrightarrow Ar^{5+} \\
KL_1L_{23} &\longrightarrow L_{23}MM + L_1L_{23}M \rightarrow (L_{23}MM + M_1M_{23})n'_p \longrightarrow Ar^{5+} \\
KL_1L_{23} &\longrightarrow L_{23}MM + L_1L_{23}M \rightarrow L_{23}MM + M_1M_{23} \longrightarrow Ar^{6+}
\end{aligned} \tag{34}$$

If we look at the ion charge-state fractions at energies above the K -threshold (3206.3eV), we see from Table 1 Section 1.2 that Ar^{1+} , Ar^{2+} cannot be produced through a K -shell Auger process. The possible radiative and non-radiative decay paths that lead to photoions with charge state Ar^{1+} , Ar^{2+} and Ar^{3+} are the following:

$$\begin{aligned}
\{K\} &\longrightarrow \{L\} + \omega_1 \rightarrow \{M\} + \omega_2 \rightarrow Ar^{1+} \\
\{K\} &\longrightarrow \{M\} + \omega_3 \rightarrow Ar^{1+} \\
\{K\} &\longrightarrow \{L\} + \omega_4 \rightarrow \{MM\} + e^- \rightarrow Ar^{2+} \\
\{K\} &\longrightarrow \{L\} + \omega_4 \rightarrow \{MMM\} + 2e^- \rightarrow Ar^{3+} \\
\{K\} &\longrightarrow \{MM\} + e^- \rightarrow Ar^{2+}
\end{aligned} \tag{35}$$

where the vacancies produced in a given shell are shown in curly brackets. From the photoion charge state fraction measurements for photon energies above 3210eV we are able to obtain the fluorescence yield for the *Ar* *K*-hole as well as the Auger yield.

2 Experimental Setup

Three distinct experiments were carried out for this thesis. We have measured the electron yield from the doubly excited states of He produced by fast carbon impact as a function of the electron emission energy and angle. We have also measured the angular correlation between the Ar *KLL* and *LMM* Auger electrons coming from the cascade decaying process after *K*-shell photoexcitation. In addition we have studied how the experimental parameters can affect the measurements of the Ar photoion charge state distribution. After gaining sufficient control over the experimental parameters, we made non-coincidence measurements of the Ar photoion charge-state distribution as a function of the photon energy, and coincidence measurements of the Ar photoion charge-state distribution with the Ar *KLL* Auger electrons. We will describe the experimental setup for all three experiments.

2.1 Doubly Excited States of Helium Produced by Fast-Ion Impact

We have measured the energies and emission angles of the electron yields from the doubly excited states $2s^2(^1S)$, $2p^2(^1D)$, $2s2p(^1P)$ and $2p^2(^1S)$ in helium produced by fast carbon-ion impact.

The C^{q+} ($q=4,5,6$) ions used in the experiment are provided by the 7.0MV EN-Tandem Van de Graaff Accelerator at the Oak Ridge National Laboratory. Ions of 1.67MeV/u energy enter the experimental chamber, interact with helium atoms in an approximately 1mm diameter diffusive gas jet crossing the ion beam and are finally collected in a Faraday cup for normalization purposes. The typical beam current in the Faraday cup is $\sim 0.9nA$. The beam is collimated by two apertures of 1mm and 0.75mm diameters, respectively, located 2m apart. To provide a differential pumping stage between the experimental chamber and the 90° bending magnet of the accelerator a, 5mm diameter differential pumping hole is placed 26cm from the

experimental chamber. The electrons emitted from the interaction region are energy and angle analyzed by a double-pass cylindrical mirror analyzer (CMA). The gas needle as well as the ion beam are positioned at the focal point of the CMA, as shown in Fig. 4.

This position is critical for obtaining the best energy resolution achievable with the CMA. After the electrons pass through the analyzer they are detected with a 40mm diameter position sensitive detector (PSD) mounted 17.52mm down stream from the CMA exit aperture. The PSD consists of a pair of microchannel plates (MCP) and a resistive anode allowing us to obtain the angular distribution of the measured electrons for each pass energy of the CMA. For more details about the PSD see Section 2.1.2 and Appendix A. Between the CMA exit aperture and the PSD there are two grids. The first grid is grounded at 13.56mm from the CMA

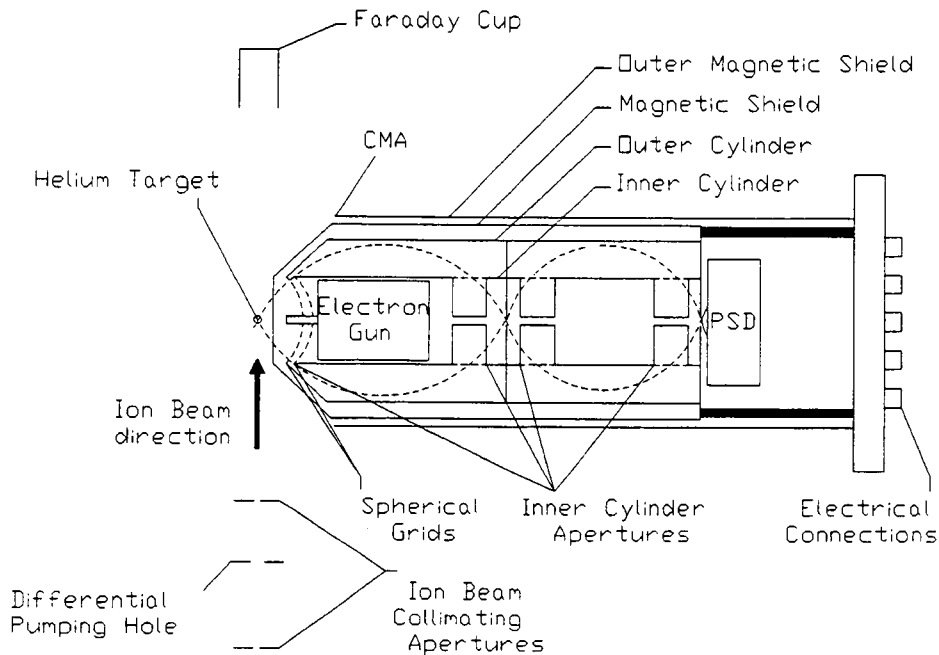


Figure 4: Schematic drawing of the apparatus used in the measurements of autoionized electrons from the doubly excited states of He, top view (not to scale).

exit aperture, and the second grid is at an intermediate voltage between ground and the voltage applied to the first surface plate of the PSD, located 0.76mm downstream from the first grid. These grids were installed to ensure that the electrons will hit the sensitive area of the PSD at an angle near to 90° with respect to the first (or entrance) plate of the PSD detector with sufficient kinetic energy for a better detection efficiency. As the MCP's first surface plate is composed of channels with a bias angle of 8° with respect to the surface normal, their detection efficiency will in general depend of the angle at which the electrons hit the MCPs. The exception is the case where the electrons hit normal to the MCP surface, in which case a more uniform detection efficiency over the entire detector surface will result. The experimental chamber is covered inside with CO-NETIC AA foil [34] that has magnetic shielding properties similar to μ -metal, so that the electron trajectories are minimally affected by external magnetic fields. The background pressure in the chamber is $2.6 \times 10^{-8} \text{mbar}$. With the high purity He gas, the chamber pressure is $8.7 \times 10^{-6} \text{mbar}$.

The CMA is ususally operated in its high-resolution mode, so the electrons pass through the analyzer with an energy of 5eV , for which the energy resolution is 0.1eV . This high resolution is accomplished by a pair of spherical grids in the entrance of the CMA. The first grid is grounded to guarantee a field-free source region while the second grid is at a potential such that it reduces the electron kinetic energy to the desired pass energy, in this case to 5eV [35]. For more details about the CMA see Section 2.1.1. The inner- and outer-cylinder voltages in the CMA are scanned together with the grid voltage so that electrons with kinetic energies between 32eV and 42eV are analyzed by the CMA. The voltage difference between the inner and outer cylinder is kept constant and is given by $V_1 = V_{IC} - V_{OC}$ such that $5\text{eV} = 1.7 \times eV_1$. The factor 1.7, the analyzer constant, relating kinetic energy and voltages, is geometrical in origin. When operating the CMA in its high resolution mode, the source size is reduced from a 4mm diameter area to a 1.5mm diameter

area, resulting in a reduction of the transmission. The CMA transmission is such that the image space is a one to one replica of the object space, with unit magnification, except for aberrations. The acceptance angle of the CMA is of $42.3^\circ \pm 3^\circ$, so we are able to measure electrons emitted between 47.7° and 132.3° with respect to the ion beam direction. The angle transformation between the real space and the image plane is given by $\theta = \cos^{-1}(\sin 42.3^\circ \sin \phi)$ where θ is the electron emission angle with respect to the beam direction in the interaction region (or real space) and ϕ is the angle measured in the detector plane (or image plane) as shown in Fig. 5. One half of the image on the PSD is identical to the second half where the ion beam direction is the symmetry axis. The inner cylinder of the CMA has four apertures through which the electrons pass, as shown in Fig. 4. These apertures are four slotted holes, separated approximately 5mm from each other, covered by 99% transmission grids. The separations between the slotted apertures produce four shadows in the image of the PSD where electrons are not detected.

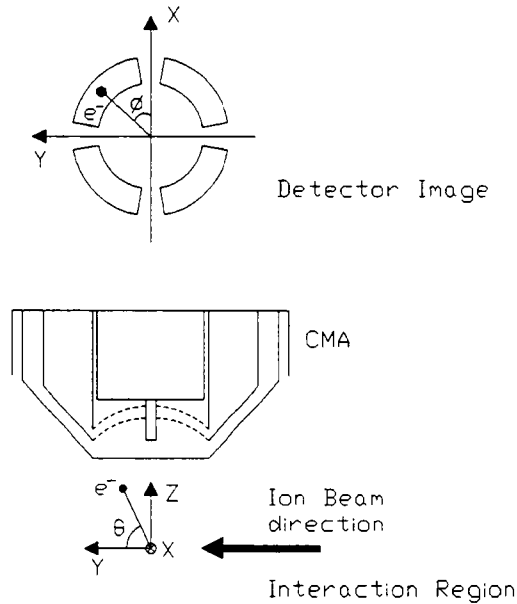


Figure 5: Schematic drawing of the angle transformation between the real and the image planes

After the electrons are detected with the PSD, the signal produced by the detector is amplified and analyzed by an acquisition system specially designed for this kind of application. The acquisition system is composed of a timer/counter, a coincidence module, an event module and a digital to analog converter (DAC). All these modules are built to the CAMAC standard so that a CAMAC crate controller is used to communicate among these devices with the aid of a personal computer (PC) data acquisition and analysis program called MP. For a more detailed description about the capabilities of the MP program see Appendix D. A schematic drawing of the acquisition system is shown in Fig. 6; a more detailed diagram is shown in Appendix A. When electrons hit the detector, the resistive anode distributes the signal to each of its four corners. The signal amplitudes carry the electron position information. These signals are amplified and analyzed by an Imaging Particle Detector, giving a digital output with the x and y position where the electron hit the detector (see Section 2.1.2 for more details).

These signals are sent to the event module that communicates to the CAMAC

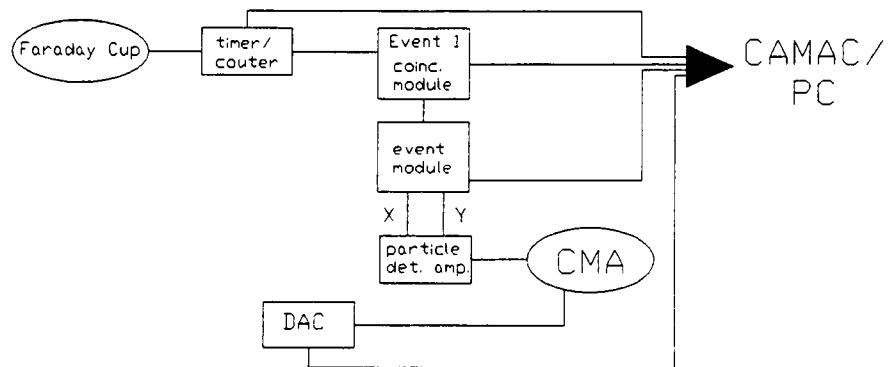


Figure 6: Schematic drawing of the electronics used in the measurements of autoionized electrons from the doubly excited states of He.

line and to the coincidence module that an event is ready to be analyzed. The coincidence module accumulates events (electrons with a given kinetic energy) for a preset number of ions hitting the Faraday Cup. The number of ions accumulated for a given electron kinetic energy is set by the timer/counter module. After a given number of ions hit the Faraday cup, the DAC (through the MP program) steps the voltages on the inner and outer cylinders of the CMA to let electrons with variable kinetic energy sequentially pass through the analyzer. In this way we are able to obtain the energy and emission angle of the electron-emission yield from the doubly excited states $2s^2(^1S)$, $2p^2(^1D)$, $2s2p(^1P)$ and $2p^2(^1S)$ in helium produced by fast carbon-ion impact.

2.1.1 Detailed Description of the CMA

The double-pass cylindrical mirror analyzer (CMA) used in these experiments, which incorporates spherical grids for preretardation, is shown in Fig. 4. The acceptance angle $\alpha = 42.3^\circ$ is the required value to satisfy second-order focusing conditions at the circular aperture between the two CMA stages, and at the exit of the second stage [35]. The trajectories of electrons which pass through the second stage must originate from a well-defined source area (of radius $\sim \rho$) when the source is positioned at the focal point (or source point) of the first stage. For an isotropic distribution within the source volume of the analyzer, the probability, T' , for an electron to pass through both stages is

$$T' = \Omega T / 4\pi \quad (36)$$

where Ω is the solid angle subtended by the CMA entrance aperture from the electron source point and T is the analyzer transmission. Electrons which originate within ρ of the analyzer source point are transmitted through both stages with a probability factor $T \approx (1 - \frac{0.01}{\sin 42.3^\circ})^4 \times (0.99)^2$ for 99% transmission grids; as the electrons pass through four grid at 42.3° and through two at 90° .

In the high resolution (or pre-retardation) mode of operation of the CMA the spherical grid nearest to the interaction region (or source point) remains at ground potential and the second grid, which is in electrical contact with the inner-cylinder assembly, is driven negative. To accurately assess the source volume in this mode of operation, it is necessary to consider refraction in the retarding-field region between the grids. As shown in Fig. 7, an electron which originates a distance, x_0 , from the analyzer source point is refracted by the retarding field such that its extrapolated trajectory is at a greater distance, x' , from the source point. To determine the relationship between x_0 and x' let's consider the case where the initial kinetic energy, E_k , of an electron is reduced to a passing energy, E_p , by the retarding field. Then one obtains that

$$\frac{E_p}{E_k} = \left(\frac{v'}{v_0} \right)^2 = \frac{v_x'^2 + v_y'^2}{v_{0x}^2 + v_{0y}^2} \approx \left(\frac{v_y'}{v_{0y}} \right)^2 \quad (37)$$

for $v_y \gg v_x$. Since the retarding field is parallel to the y axis, $v_{0x} = v_x'$ and

$$\frac{\theta_0}{\theta'} = \frac{x_0}{x'} = \frac{v_y'}{v_{0y}} = \left(\frac{E_p}{E_k} \right)^{\frac{1}{2}} \quad (38)$$

or

$$x' = x_0 \left(\frac{E_k}{E_p} \right)^{\frac{1}{2}} \quad (39)$$

For a retarded electron to pass through the double-pass CMA, it is clear that the virtual source point, x' , must lie within the sphere of radius, ρ , which means that the real source point x must lie within a sphere of radius

$$\rho' = \rho \left(\frac{E_p}{E_k} \right)^{\frac{1}{2}} \quad (40)$$

Thus, the effect of retardation is to improve the energy resolution by the ratio E_k/E_p and to collapse the source volume according to equation (40).

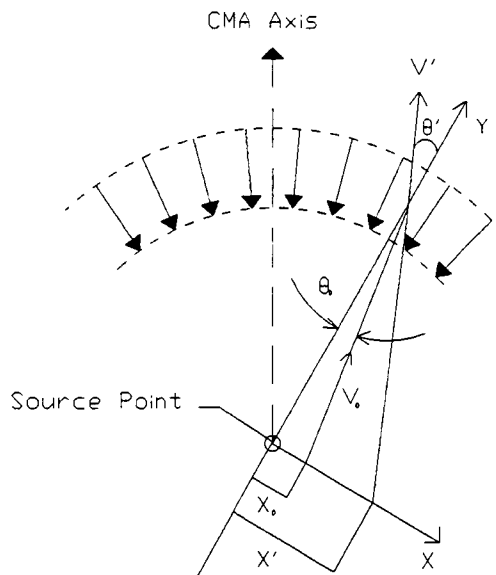


Figure 7: Schematic diagram illustrating refraction by retarding grids.

In addition to refraction, the grids also affect the analyzer operation in the retarding mode because of grid scattering. When an electric field terminates on a grid of finite mesh, the equal-potential surfaces are rippled close to the grid. Each grid opening acts as a lens which causes the transmitted electrons to be scattered from their initial trajectory. Grid scattering reduces the analyzer transmission, T , and increases the imaged area, A , such that the luminosity, $T \times A$ is not affected. In our case where the source region is small compared to the imaged area, the effect of the grid scattering is to reduce the transmission by about a factor of E_k/E_p . The CMA resolution $R = \Delta E/E_p$ is 1.6% without retardation. The specifications for the CMA are shown in Table 2.

For the CMA to work in these two possible modes (retarding and nonretarding), a programmable power supply (model PHI 20-805) determines the voltage applied

Table 2: CMA specifications.

Analyzer Energy Resolution	
Without Retardation	
Focused Source	1.6%
Flooded Source	2.0%
With Retardation	$0.02E_p$
Source Size	
Without Retardation	4mm diam. area
With Retardation	
Theoretical	$4.0(E_p/E_k)^{\frac{1}{2}}mm$ diam. area
Experimental	< 4mm diam. area for $E_k/E_p > 40$
Analyzer Constant	1.7
CMA Transmission (approx.)	92%

to the inner and outer cylinders. In the nonretarding mode, the inner cylinder is always at the ground potential, while the outer cylinder has a variable voltage (V) applied such that electron energy = $1.7 \times V$, where 1.7 is the analyzer constant. In the high resolution mode the power supply works as shown in Fig. 8. In this mode of operation the voltages applied to the inner and outer cylinders vary according to the variable supply voltage, V_3 , but keeping a constant voltage difference between them, V_1 . The pass energy is determined by the voltage V_1 . The offset supply, V_2 , automatically tracks the pass-energy supply, V_1 , according to the formula $V_2 = 1.7 \times V_1$. In this retarding mode of operation the electron energy is precisely equal to eV_3 . In this way the desired energy resolution can be chosen by adjusting V_1 without affecting the measured electron energy.

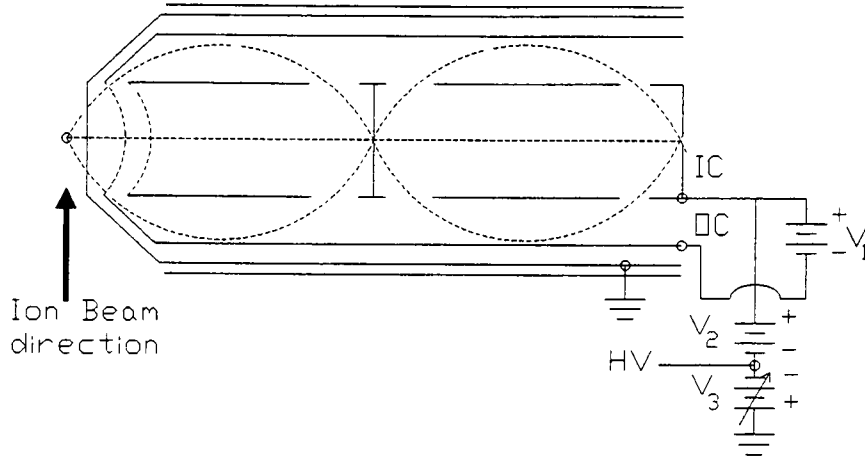


Figure 8: Schematic diagram of the CMA power supply working in retarding mode.

Summarizing the relations of the voltages applied to the CMA when working in high resolution mode,

$$\begin{aligned} V_1 &= V_{IC} - V_{OC} \\ PassEnergy &= 1.7 \times eV_1 \\ V_2 &= 1.7 \times V_1 \end{aligned} \quad (41)$$

$$ElectronEnergy = eV_3$$

In our case (when working in the high resolution mode), the voltages applied to the inner and outer cylinders are scanned with the MP program. Through the program, we can set the desired energy range by setting the starting voltage in the DAC (V_{dac} Start) and the stopping voltage in the DAC (V_{dac} Stop). The CMA power supply requires 0 – 10V (programmable DC voltage) at its input to drive V_3 .

Table 3 shows the voltages applied at the power supply input, and the measured voltages at the inner and outer cylinders and the value of V_3 (that is the HV output of the power supply) for different DAC voltages (V_{dac}) set with the MP program. From this table we can obtain a rough calibration of the DAC voltage set with the MP program (V_{dac}) as a function of the electron energy for each mode of operation:

$$\text{Nonretarding Mode (Auger): } V_{dac} = 0.147 \times E_e - 0.139$$

Table 3: Inner and outer cylinder voltages for different DAC voltages set with the MP program, for various modes of operation (all numerical values are given in Volts). The retarding modes used correspond to ESCA and UPS modes.

Vdac	Input	Nonretarding $V_3(\text{Auger})$	10eV Pass Energy (ESCA)			5eV Pass Energy (UPS)		
			IC	OC	V_3	IC	OC	V_3
0	0.0	-0.50	9.43	3.53	-0.50	4.950	1.999	-0.07
10	0.2	-40.5	-30.53	-36.4	-40.5	0.958	-1.99	-4.07
20	0.4	-80.5	-70.50	-76.5	-80.5	-3.05	-6.00	-8.07
30	0.6	-120.5	-110.5	-116.4	-120.5	-7.04	-10.00	-12.07
40	0.8	-160.5	-150.4	-156.4	-160.5	-11.04	-13.99	-16.07
50	1.0	-200.5	-190.6	-196.4	-200.5	-15.04	-18.00	-20.07
60	1.2	-240.5	-230.4	-236.4	-240.5	-19.04	-22.00	-24.07
70	1.4	-280.5	-270.6	-276.4	-280.6	-23.04	-26.00	-28.07
80	1.6	-320.5	-310.6	-316.4	-320.5	-27.04	-30.00	-32.07
90	1.8	-360.5	-350.0	-358.0	-360.5	-31.04	-34.00	-36.07
100	2.0	-400.5	-392.0	-396.0	-400.5	-35.04	-38.00	-40.07

Retarding Mode (ESCA): $V_{dac} = 0.250 \times E_e - 0.139$

Retarding Mode (UPS): $V_{dac} = 2.499 \times E_e - 0.132$.

In ESCA (Electron Spectroscopy for Chemical Analysis) mode $0 \leq V_3 \leq 2000V$, and in the UPS (Ultraviolet Photoelectron Spectroscopy) mode $0 \leq V_3 \leq 200V$ For more accurate energy calibrations for different electron energy regions see Appendix B.

2.1.2 Description of the Position Sensitive Detector

As discussed above, electrons analyzed by the CMA were measured by a position sensitive detector (PSD). This detector is composed of a resistive anode placed behind two 40mm diameter microchannel plates, so each incident electron produces a detectable charge pulse on the anode. The resistive anode has four output contacts (denoted A, B, C and D) in its respective corners and produces a charge pulse in each output related to the position of the charge pulse produced by the detected electron. The position analyzer electronics generates a voltage output proportional to the pulse-position coordinates of the form (details about the connection are shown in Appendix A)

$$\begin{aligned} X &= \frac{Q_B + Q_C}{Q_A + Q_B + Q_C + Q_D} \\ Y &= \frac{Q_A + Q_B}{Q_A + Q_B + Q_C + Q_D} \end{aligned} \quad (42)$$

where Q_A , Q_B , Q_C and Q_D denote the charge collected at each corner. Once a pulse has been analyzed and its position signal is available at the X and Y outputs, a Strobe pulse (and a Z -axis TTL pulse) is sent. This occurs $18\mu\text{sec}$ after the electron hits the detector for the analog signals and $23\mu\text{sec}$ for the digital signals. This detector system also has a shaper amplifier that converts the low-level charge pulse into a higher-level bipolar pulse (Sum signal) that is used in the coincidence measurements. This signal comes $2.75\mu\text{sec}$ after the electron hits the detector and is $4\mu\text{sec}$ long. A valid conversion requires the sum signal mean height to be between 1 and 1.5V; this is adjusted with the high voltage applied across each microchannel plate, being careful not to exceed 1000V per plate. The dark counts (or background flux) of the detector used is less than 5counts/sec over the entire image.

2.2 Cascade-Decay Processes Produced in Photon-Atom Interactions

We have measured the angular correlation between two Auger electrons coming from a cascade-decay process after photoexcitation. We performed electron-electron coincidence measurements between KLL and LMM Auger electrons emitted from Ar after $1s \rightarrow 4p$ photoexcitation. The experimental setup is similar to the one described in Section 2.1. For this experiment we used beam line X24A at National Synchrotron Light Source (NSLS) at Brookhaven National Laboratory. The beam line was operating with a Ge(111) double-crystal monochromator allowing us to obtain an energy range from $3000eV$ to $5200eV$, with a photon flux of $10^{12}photons/sec$ at $100mA$ ring current and $2.5GeV$ electrons energy with a resolving power $E/\Delta E = 3800 - 4500$ [36]. At this beam line the photon beam is nearly linearly polarized with the electric vector pointing in the plane of the ring. The photon beam energy was calibrated by measuring an argon absorption scan in the vicinity of the $1s \rightarrow 4p$ resonance peak at $3203.54eV$, as shown in Fig. 9 [41]. The signals for the absorption scan are the KLL Ar Auger electrons measured by an electron time-of-flight (TOF) spectrometer, discussed below.

We used an experimental chamber covered on the inside with a CO-NETIC AA foil (similar to μ -metal) so that the electrons are minimally affected by external magnetic fields. The background pressure in the chamber was $2 \times 10^{-8}mbar$. With the high-purity Ar gas the chamber pressure is $1.4 \times 10^{-5}mbar$. In order to work at this pressure we used a $10\mu m$ thick Beryllium window between the chamber and the beam line, to maintain beam line pressure of $1 \times 10^{-10}mbar$. The photon beam was collimated by a $1.5mm$ diameter aperture placed $33cm$ from the interaction region. The photon beam interacted with an Ar effusive gas target. Electrons sequentially emitted from the interaction region were energy and angle analyzed by two analyzers: (1) by an electron time-of-flight analyzer (TOF) that views the interaction region in

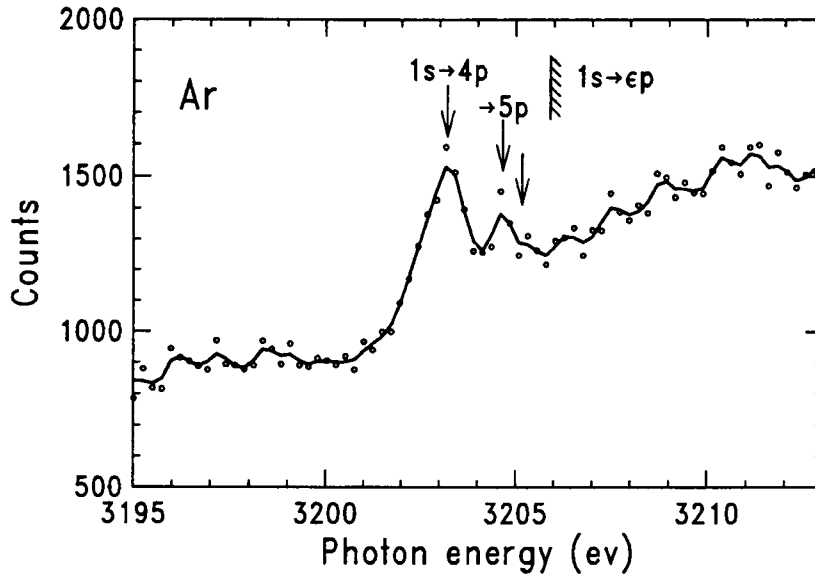


Figure 9: Absorption scan to locate the $1s \rightarrow 4p$ resonance in argon.

the photon beam polarization direction; and (2) by the CMA mounted perpendicular to the electron TOF axis and to the beam orientation as shown in Fig. 10.

Electrons emitted between 47° and 132° with respect to the photon beam direction are measured by the CMA. In this case the image obtained on the PSD is symmetric with respect to the photon beam polarization direction. The gas target, an effusive gas jet, was obtained from a hypodermic needle placed at 45° with respect to the photon beam, and pointing away from the electron TOF and CMA axis. The photon beam flux was monitored by the photon current produced in a stainless-steel plate, measured by an electrometer and sent to a current digitizer, and used for normalization purposes. The TOF analyzer entrance aperture is 28mm away from the interaction region. Its flight path is pumped by a 70l/s turbo molecular pump to avoid inelastic scattering of electrons at the gas in the TOF analyzer. The flight path is shielded by μ -metal from magnetic fields which may cause a loss of transmission of the analyzer. The electrons detected by the electron TOF microchannel plate (MCP) detector produce a signal used as the start pulse in the time-to-digital converter (TDC) electronic module. In this case the stop pulse is the bunch marker

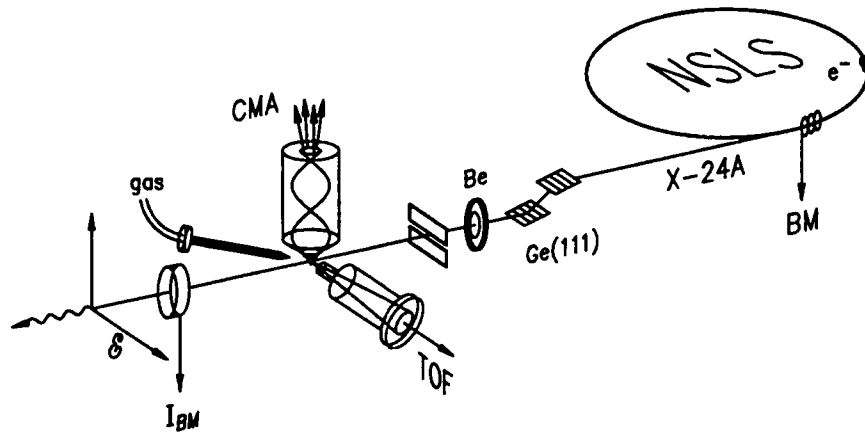


Figure 10: Schematic drawing of the experimental setup for the electron-electron coincidence measurements.

(BM) signal from the storage ring (with a period of $567nsec$), gated with the signal of the electrons hitting the CMA detector. The CMA voltages are set so that different groups of *KLL* Auger electrons are selected in coincidence with the electrons measured by the electron TOF. A diagram of the electronics used is shown in Fig. 11. For more details about the electronics see Fig. 54 in Appendix A.

Both the CMA and the TOF need to be precisely aligned with respect to the gas needle and with respect to each other to be certain that they view the same interaction region. The gas needle is mounted on an XYZ translator that allows us to move and position the needle with a precision of $0.001mm$. The CMA is mounted on a $10inch$ bellows that allow us to align it with a precision of $0.16mm$. The TOF is aligned horizontally and vertically through two linear feedthroughs that allow us to align it with a precision of $0.05mm$. The electron TOF has the property that its transmission depends on the alignment, the mode of operation (retarding or accelerating mode) and on the electron energy. After the TOF was precisely aligned we obtained a transmission curve. For this purpose we recorded the photoelectron

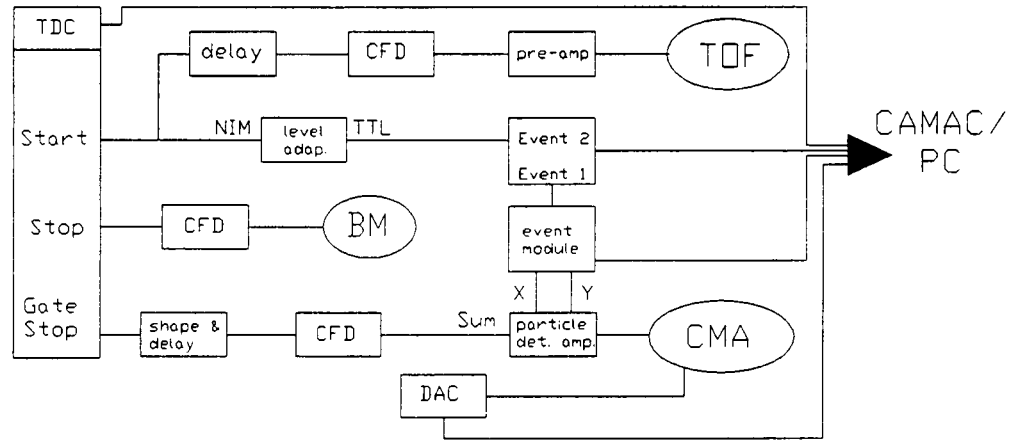


Figure 11: Schematic drawing of the electronics used for the electron-electron coincidence experiments.

energy spectra for different photon energies and observed the $1s$ photoline moving in kinetic energy (or, equivalent, time of flight) and changing its relative area with respect to the argon KLL Auger lines as shown in Fig. 12. From these spectra we obtained the transmission curve that is given by the ratio between the area of the $1s$ photoline and the area of the argon KLL Auger lines as shown in Fig. 13, this is true for kinetic energies below the $1s$ satellites threshold ($25eV$ above the $1s$ threshold). Above the satellite threshold, the satellite lines are added to the $1s$ photoline area. Experience shows that is safe to assume that the transmission curve is flat above $25eV$ kinetic energy. In this way we were able to normalize the $1s$ photoline relative to the KLL Auger lines (as the area of the KLL Auger lines do not change with photon energy). Then the ratio between the areas will only depend on the $1s$ photoelectron transmission through the analyzer for different kinetic energies. The position of the $1s$ photoline for different photon energies was used to obtain an energy calibration for the TOF spectra as shown in Fig. 14. The energy calibration curve is used to transform the time spectra obtained with the electron TOF into energy spectra, as will be shown in Section 4.2.

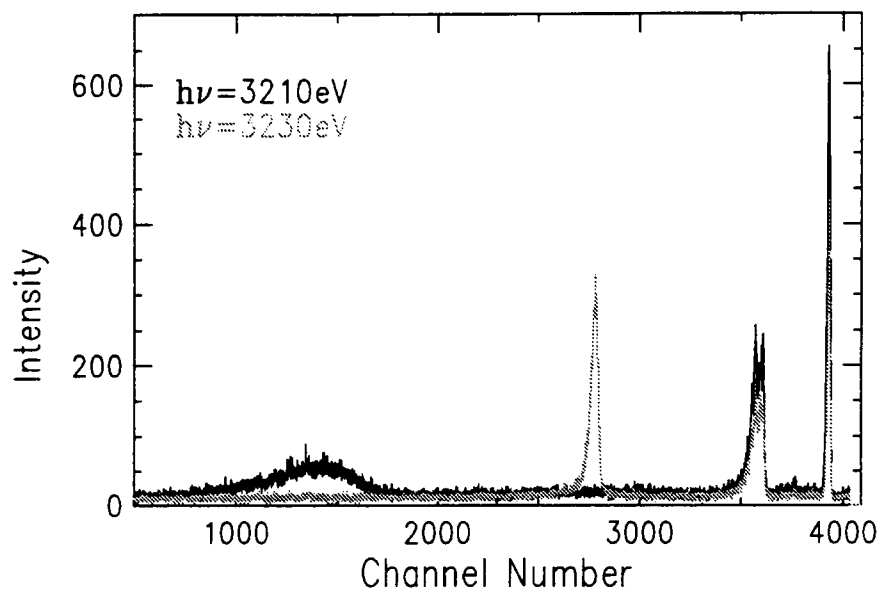


Figure 12: TOF spectra for two different photon energies.

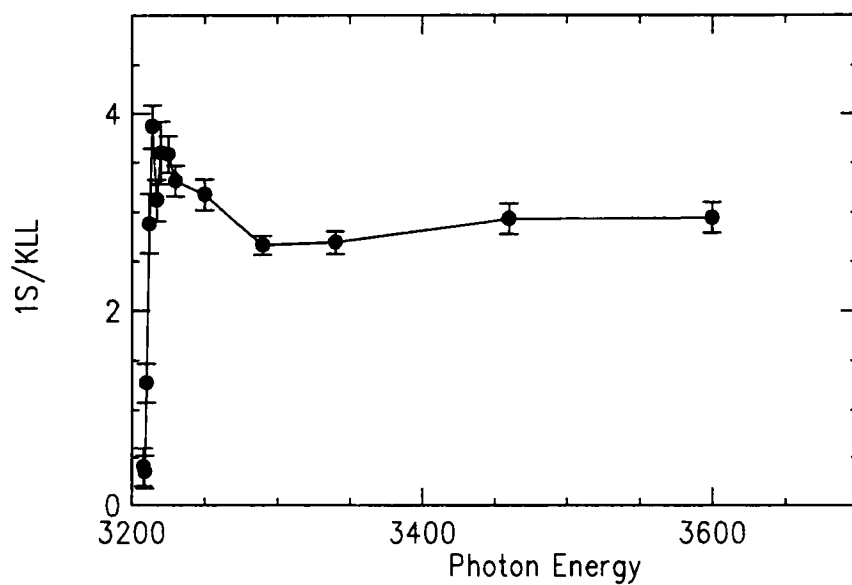


Figure 13: Transmission curve for the electron TOF with 4V acceleration.

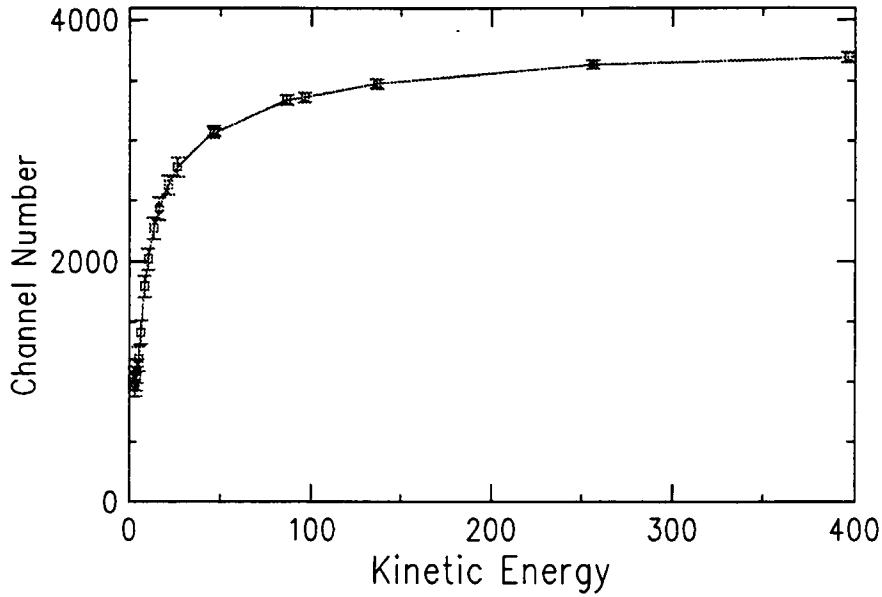


Figure 14: Energy calibration for the TOF spectra with 4V acceleration.

As was mentioned before, the electron TOF analyzer has two modes of operation, the retarding and the accelerating mode. When working in the retarding mode a negative voltage is applied to the lens system, (see Fig. 15). The electrons are slowed down during their flight through the analyzer to achieve a higher energy resolution. As the TOF has a fixed kinetic energy resolution of 2%, slower electrons permit better energy resolution. To measure very-slow electrons, the electron TOF requires operation in the acceleration mode. In this case a positive voltage is applied to the lens system so electrons with small kinetic energies can be detected. The total energy resolution of the electron TOF analyzer [42, 43] is given by

$$\frac{\Delta E}{E} = \left[\left(\frac{2\Delta t}{t} \right)^2 + \left(\frac{2\Delta z}{z} \right)^2 + \left(\frac{2\Delta \lambda}{\lambda} \right)^2 \right]^{1/2} \quad (43)$$

where E is the kinetic energy, ΔE is the energy resolution, t is the time of flight, Δt is the uncertainty in the time measurement, z is the flight-path length for the electron, Δz is the uncertainty in that length, λ is the x-ray wavelength and $\Delta \lambda$ is

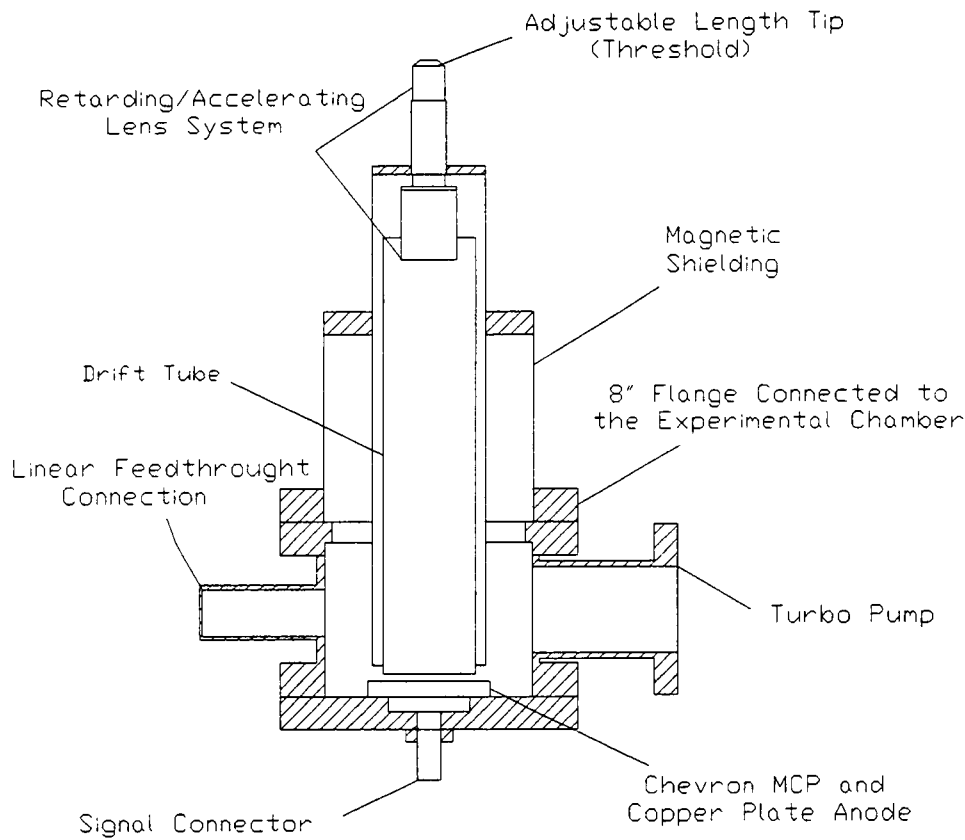


Figure 15: Schematic drawing of the electron time-of-flight analyzer.

its uncertainty.

The electron TOF analyzer used in these experiments has a total flight path of 460mm (this is the distance from the interaction point to the microchannel plate (MCP) detector), an angular acceptance of 2.6° half angle, and an energy resolution of 2%. The microchannel plate (MCP) detector is composed of two microchannel plates in Chevron operation [50] and a copper-plate anode. For a better electron-detection efficiency the two microchannel plates in the detector are insulated from each other, and a voltage is applied to both central faces of the plates such that there is a voltage difference of 200 to 400V between the centers of the MCP, as shown in Fig. 16. A voltage between 800 to 1000V is set across each MCP and the voltage

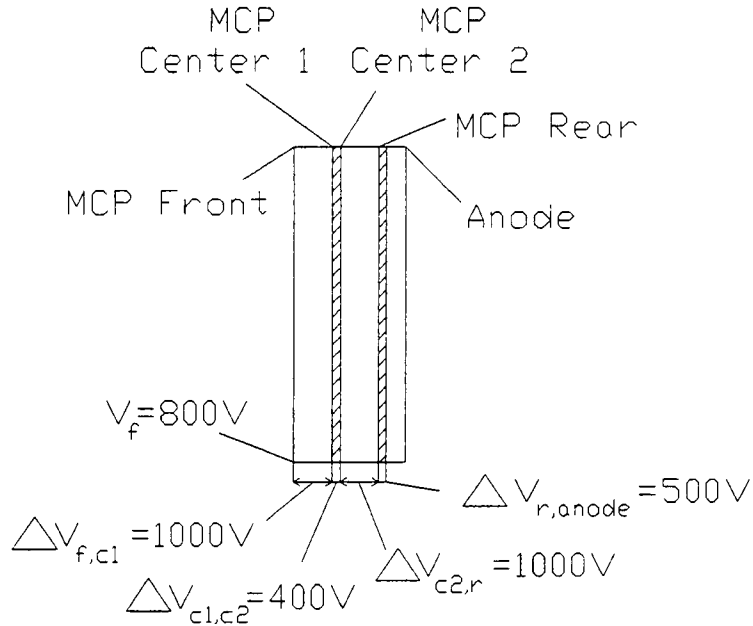


Figure 16: Schematic drawing of the microchannel plate detector.

applied to the anode is such that there is a voltage difference of approximately 300 to 500V between the anode and the rear of the MPC. The front surface of the MCP detector is held at a voltage between 600 and 800V with respect to ground, so electrons hit the MPC detector with sufficient kinetic energy for maximum detection efficiency.

2.3 Photoion Charge-State Distribution and Electron-Ion Coincidences

An experimental setup similar to the one described in Section 2.2 was used to measure photoion charge state distributions and photoions in coincidence with *KLL* Auger electrons. These experiments were also done on beam line X24A at NSLS. For these experiments the electron TOF spectrometer was removed and an ion time-of-flight (ion TOF) spectrometer was mounted in the experimental chamber, see Fig. 17 for

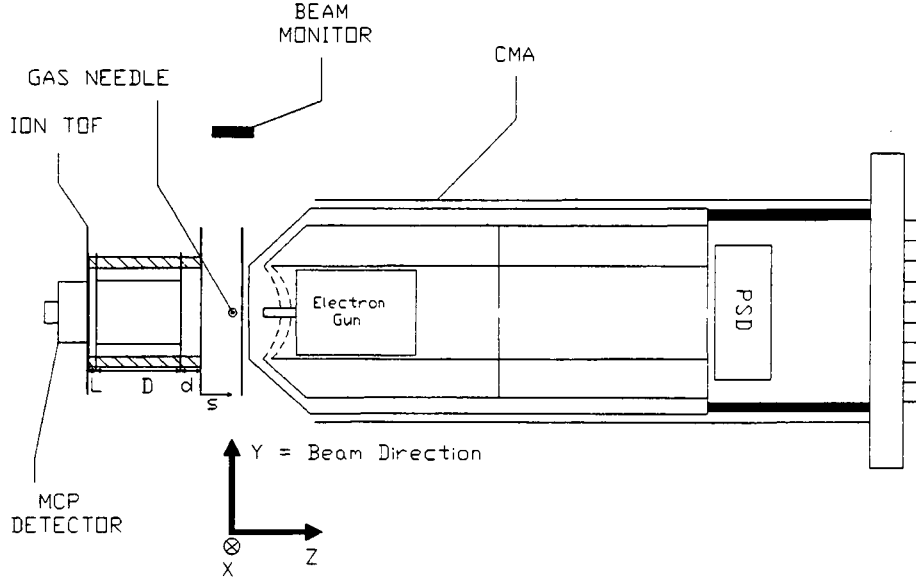


Figure 17: Schematic diagram of the experimental setup for the electron-ion coincidence measurements.

the schematic diagram of the experimental setup.

In this experimental setup both the CMA and the ion TOF spectrometers were mounted perpendicular to the photon beam and the polarization vector direction. The effusive gas source emerges from a hypodermic needle attached to the ion TOF and placed at a distance s from the extraction plate. The ion TOF was mounted on an XYZ translator, allowing us to align it with respect to the photon beam and the CMA. The gas needle was placed adjacent to the photon beam and close to the focal point of the CMA. The photon beam flux was monitored by a beam monitor. The electron current produced in the beam monitor was measured by an electrometer and sent to a current digitizer for normalization purposes. Two different setups for the electronics were used depending on whether we were requiring coincidences with the cascade *KLL* Auger electrons. The simplest setup is for the singles (non-coincidence measurements). In this case the signal, coming from the MCP detector of the ion

TOF, is the start signal of a time to amplitude converter (TAC) and the stop signal is the bunch marker (BM) signal from the storage ring, see Fig. 18. For more details about the electronics setup see Fig. 56 in Appendix A.

For the coincidence spectra, the signal coming from the MCP detector of the ion TOF is the start signal of the time-to-digital converter (TDC) and the stop signal is the signal coming from the PSD of the CMA. In this case the CMA voltages were scanned over a small energy window of a given *KLL* Auger line. The number and the size of the steps in the scan were set by a digital to analog converter (DAC) controlled by the acquisition system through the CAMAC crate controller, as shown in Fig. 19. With this setup we were also able to obtain the angular distribution of the *KLL* Auger lines for each photoion charge state.

The ion TOF that we built for these experiments, Fig. 20, is composed of a pusher plate, an extraction plate, a first drift tube, an acceleration plate and a second drift region. The first acceleration region is the one between the pusher and the extraction plates. This is the ionization (or interaction) region where the ions are formed at a distance s from the extraction plate. The region between the extraction

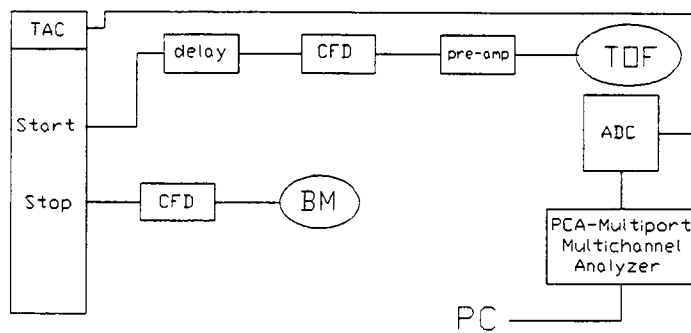


Figure 18: Schematic diagram of the electronics for the non-coincidence measurements of photoion charge state distributions.

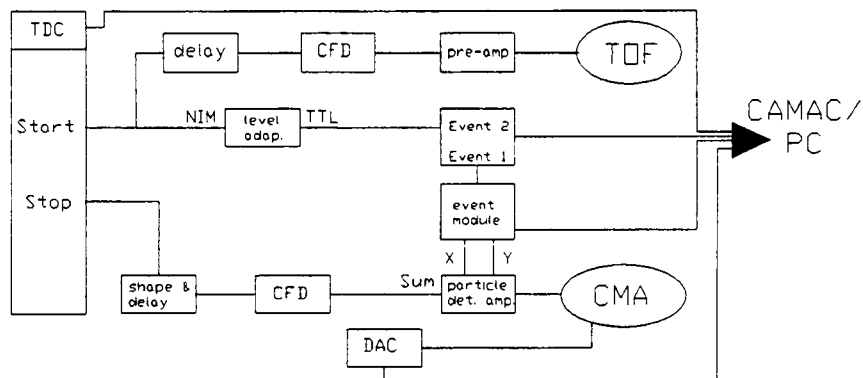


Figure 19: Schematic diagram of the electronics for the electron-ion coincidence measurements.

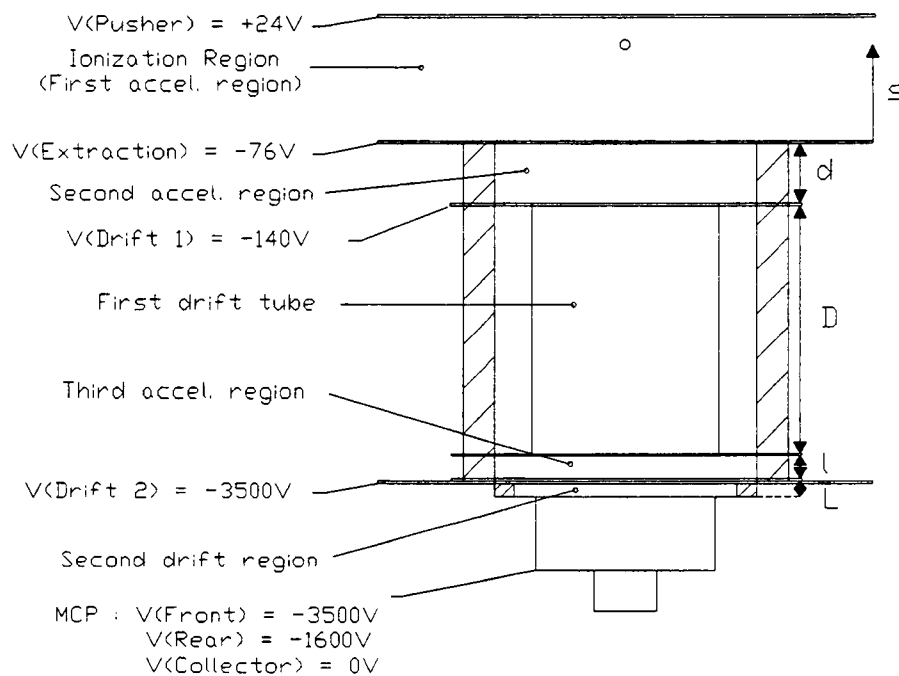


Figure 20: Schematic diagram of the ion TOF.

and the first drift region (d region) is the second acceleration region. The two drift regions (D and L regions) are field free regions (constant voltage) as shown in Fig. 20. There is also a third acceleration region (l region) that is very small. Its only purpose is to give the ions sufficient kinetic energy as they hit the MCP, because the detection efficiency of the MCP depends on the energy of the detected ions. The distance between the pusher and the extraction plate is 1.27cm . The distance between the pusher and the interaction region is $s = 9.68\text{mm}$, while the distance between the extraction plate and the entrance to the first drift region (or second acceleration region) is $d = 6.22\text{mm}$. The length of the first drift region is $D = 2.53\text{cm}$. The lengths of the third acceleration stage and the second drift region are equal to $l = 2.49\text{mm}$ and $L = 2.44\text{mm}$, respectively. The apertures of the plates are covered by a 99.8% transmission grid. The size of the apertures for each stage are the following: pusher= 2.53cm , extraction= 4.72mm , drift 1= 1.27cm , 1.91cm and the drift 2= 2.54cm . The MCPs used are non-imaging-quality, and have 18mm diameter active area. The total length of the ion TOF including the MCP is approximately 10cm . The dimensions of the ion TOF spectrometer were chosen such that we have good time and space resolution as will be discussed later in Section 2.3.1 [44]. When the ions are created in the ionization region they will have an initial energy U_0 (given by the thermal energy). They will then increase their energy as they fly through the ion TOF to a value U_f given by

$$U_f = U_0 + qsE_s + qdE_d + qlE_l \quad (44)$$

where q is the ion charge state, in units of e (the electron charge), and E_s , E_d , E_l are the electric fields in the s , d and l regions respectively. U_f is independent of the mass m of the ion. The total flight time is:

$$T_t(U_0, s) = T_s + T_d + T_D + T_l + T_L \quad (45)$$

The last two terms in this expression are very small compared to the first three due

to the short l and L distances and the high voltages applied across the l stage. If we restrict our analysis to a one-dimensional problem the time-of-flight in each region is given by (in a three dimensional problem U_0 must be replaced by $U'_0 = \frac{1}{2}mv_{0||}^2$, where $v_{0||}$ is the initial ion velocity in the collector direction)

$$\begin{aligned}
T_s &= \frac{(2m)^{1/2}}{qE_s} \left\{ (U_0 + qsE_s)^{1/2} \pm (U_0)^{1/2} \right\} \\
T_d &= \frac{(2m)^{1/2}}{qE_d} \left\{ U^{1/2} - (U_0 + qsE_s)^{1/2} \right\} \\
T_D &= (2m)^{1/2} D/2U^{1/2} \\
T_l &= \frac{(2m)^{1/2}}{qE_l} \left\{ U_f^{1/2} - U^{1/2} \right\} \\
T_L &= (2m)^{1/2} L/2U_f^{1/2}
\end{aligned} \tag{46}$$

The $+$ and $-$ signs in T_s corresponds to initial velocities directed away from and towards the collector, respectively [44], and $U = U_0 + qsE_s + qdE_d$.

For these experiments two different set of voltages were used in the ion TOF. For the non-coincidence measurements we used relatively high voltages in the ion TOF spectrometer to assure that all charge-state fractions from Ar^{1+} to Ar^{7+} were in a time range smaller than $567nsec$ (equal to the time difference between two consecutive photon pulses from the storage ring). The voltages used were: $V(Pusher) = 240V$, $V(Extraction) = -760V$, $V(Drift1) = -1200V$, $V(Drift2) = -4350V$ and $V(MCPfront) = -4350V$. For the coincidence measurements, the voltages were smaller to avoid materially affecting the trajectories of the electrons to be detected in the CMA. In this case the voltages used were: $V(Pusher) = 24V$, $V(Extraction) = -76V$, $V(Drift1) = -120V$, $V(Drift2) = -4350V$ and $V(MCPfront) = -4350V$. For these two set of voltages, typical ion TOF parameters are shown in Table 4, where the first set of values corresponds to the smallest set of voltages and the second set of values corresponds to the highest set of voltages in the ion TOF, q is the ion charge state ($q = 1...7$) and the initial energy is approximated by $U_0 \approx 0$. The energies

Table 4: Typical parameter values for the ion TOF analyzer.

Distance (cm)	Electric field (volts/cm)	Energies (eV)	Time of Flight (μsec)
$s = 0.968$	$E_s = 78.74$	$U_s = q76.12$	$1.009/\sqrt{q}$
$d = 0.622$	$E_d = 70.74$	$U_d = q120.06$	$0.287/\sqrt{q}$
$D = 2.53$	$E_D = 0$	$U_D = U_d$	$1.051/\sqrt{q}$
$l = 0.249$	$E_l = 16987.95$	$U_l = q4344.6$	$0.029/\sqrt{q}$
$L = 0.244$	$E_L = 0$	$U_L = U_l$	$0.017/\sqrt{q}$
$s = 0.968$	$E_s = 787.40$	$U_s = q761.2$	$0.319/\sqrt{q}$
$d = 0.622$	$E_d = 707.40$	$U_d = q1200.6$	$0.091/\sqrt{q}$
$D = 2.53$	$E_D = 0$	$U_D = U_d$	$0.332/\sqrt{q}$
$l = 0.249$	$E_l = 12650.6$	$U_l = q4346.6$	$0.022/\sqrt{q}$
$L = 0.244$	$E_L = 0$	$U_L = U_l$	$0.017/\sqrt{q}$

defined in Table 4 correspond to the kinetic energy at the end of each region; the electric fields are assumed to be approximately uniform between each pair of plates i and $i + 1$, and are given by $E_i = (V_i - V_{i+1})/R$, where R is the distance between plates.

2.3.1 Space and Time Resolution of the Ion TOF Spectrometer

The voltages and distances between the plates of the ion TOF were chosen such that the all measurable charge states q are resolved for a particular atom with mass m . The values for the voltages and distances are determined by the space and energy resolution of the ion TOF. To study the space and energy resolution we shall assume that $U_0 = 0$, $s = s_0$ and define

$$\begin{aligned}
U_t &= qs_0E_s + qdE_d \\
k_0 &= (s_0E_s + dE_d)/s_0E_s
\end{aligned} \tag{47}$$

and

$$T(U_0, s) = T_s + T_d + T_D \tag{48}$$

where T_s , T_d and T_D are defined in equation (46) in Section 2.3. Then we can write $T(0, s_0)$ as

$$T(0, s_0) = \sqrt{\frac{m}{2U_t}} \left(2s_0k_0^{\frac{1}{2}} + d\frac{2k_0^{\frac{1}{2}}}{k_0^{\frac{1}{2}} + 1} + D \right). \tag{49}$$

Space focusing depends on the fact that an ion initially closer to the detector (and thus of smaller s value) acquires less energy and is therefore eventually overtaken by ions which have larger initial s values [44]. To find the position at which ions whose initial s values were $s = s_0 \pm \frac{1}{2}\delta s$ pass each other, we set $(dT/ds)_{0,s_0} = 0$. Using equation (48) for T and by substitution from equation (47) we obtain

$$D = 2s_0k_0^{\frac{3}{2}} \left(1 - \frac{d}{s_0} \frac{1}{k_0 + k_0^{\frac{1}{2}}} \right). \tag{50}$$

This focus condition is the same for all ions and is independent of the total energy of the system. If s_0 , d and D are fixed, E_d/E_s is uniquely determined by equation (47), since k_0 can have only one physically significant value. Then the space focus is obtained by adjusting E_d/E_s . The focus condition, $(dT/ds)_{0,s_0} = 0$, indicates that $T(0, s)$ has either a maximum, minimum or point of inflection at $s = s_0$. The point of inflection occurs when $(d^2T/ds^2)_{0,s_0} = 0$ which requires in addition to the focus condition that

$$\frac{d}{s_0} = \frac{D}{2s_0} \left(\frac{k_0 - 3}{k_0} \right). \tag{51}$$

If d/s_0 is larger than this value, $T(0, s_0)$ is a minimum point; if d/s_0 is smaller, then $T(0, s_0)$ is a maximum point. The parameter values are usually chosen such that $T(0, s_0)$ is a maximum for a better resolution.

By utilizing the series expansion of $T(0, s)$ about s_0 , we may show that $\Delta T_{\Delta s}$, the change in flight time corresponding to a small change Δs in s about s_0 , is given by

$$\Delta T_{\Delta s} = \sum_{n=1}^{\infty} \left(\frac{d^n T}{ds^n}(0, s) \right)_{s_0} (\Delta s)^n. \quad (52)$$

The measure of space resolution, M_s , is the maximum value m for which $\Delta T_{\Delta s} \leq T_{m+1} - T_m$. The time separation of adjacent masses is obtained from equation (46) or (49),

$$T_{m+1} - T_m = \left[\left(1 + \frac{1}{m} \right)^{\frac{1}{2}} - 1 \right] T_m \approx \frac{T_m}{2m}. \quad (53)$$

The first term in equation (52) is zero by the focusing condition (equation (50)). If $T(0, s_0)$ is either a maximum or a minimum point, $\Delta T_{\Delta s}$ can be evaluated from equation (52) with Δs replaced by the maximum deviation of s from s_0 , $\Delta s/2$. The successive terms are complicated but can be simplified by the use of approximations without introducing appreciable errors. Only the first non zero term of the series will be used. If $k_0 \gg 1$ and $k_0 \gg d/s_0$ equations (49), (52) and (53) give

$$M_s \approx 16k_0 (s_0/\Delta s)^2. \quad (54)$$

Equation (54) shows that the space resolution can be made adequate if s_0 is somewhat larger than Δs . This is dependent on D being large since D/s_0 determines k_0 from equation (50). Increasing d (as this increases k_0) improves the space resolution. As d/s_0 becomes larger (approaching k_0) the value of M_s is considerably larger than that given by equation (54).

To analyze the energy resolution we need to investigate the effect of the initial

velocities v_0 that will lead to an initial energy $U_0 = \frac{1}{2}mv_0^2$, where m is the mass. Assume that two ions are formed in the same initial position s , with equal speeds but in opposite directions. Let ion 1 be initially directed towards the detector and 2 initially directed away from the detector. Ion 2 will first move away from the extraction plate, decelerating due to E_s until it stops. It is then accelerated, returning to s with its original speed. Subsequently, its motion is identical with that of ion 1 which it continues to lag by the "turnaround", ΔT_θ , (i.e. twice the deceleration time). Assuming that the ion has the maximum initial energy,

$$\Delta T_\theta = \frac{2v_0m}{qE_s} = \frac{2(2mU_0)^{\frac{1}{2}}}{qE_s}. \quad (55)$$

If we consider only the effects of the initial energy, the maximum resolvable mass, M_θ , is that for which the energy time spread, ΔT_θ , equals the time between adjacent mass peaks, $T/2m$ (i.e. $M_\theta = T/2\Delta T_\theta$). This is evaluated by substituting from equation (49) and (55). If D/s_0 from equation (50) is substituted in the result, M_θ is given by

$$M_\theta = \frac{1}{4} \left(\frac{U_t}{U_0} \right) \left(\frac{k_0 + 1}{k_0^{\frac{1}{2}}} - \left(\frac{d}{s_0} \right) \frac{k_0^{\frac{1}{2}} - 1}{k_0 + k_0^{\frac{1}{2}}} \right) \quad (56)$$

M_θ increase with increasing D and decreasing s_0 and d .

The maximum resolvable mass, M , depends on the initial space and energy distribution functions. Although the exact problem is very difficult to solve, we can give the upper and lower bounds on the value of M from the previous discussion. M will be smaller than either M_s or M_θ . On the other hand it will be at least as large as the value, $M_{s,\theta}$, obtained by assuming that the total time spread is the sum of the energy and space time spreads,

$$\frac{1}{M_{s,\theta}} = \frac{1}{M_s} + \frac{1}{M_\theta}. \quad (57)$$

In our case we are looking at different charge states of the same argon ions (Ar^{q+}),

so we are interested in knowing the change in flight time for each charge state due to small changes Δs about s_0 , and to effects of the initial energy of the ions. We can estimate $\Delta s \approx 0.5mm$ because the size of the interaction region in our case is given by the intersection of the ion beam and the diffusive gas jet that is approximately $1mm^2$. To estimate the initial energy we assume that all motion is due to room temperature thermal energy. Then

$$U_0 = \frac{n}{2}kT, \quad (58)$$

where n is the number of degrees of freedom ($n = 3$), k is Boltzmann's constant and T is the temperature in Kelvin. Then at room temperature $T = 20^\circ C = 293.1K$ and $U_0 = 0.038eV$. We can estimate the change in flight time from equation (55) and (52) (keeping only the $n = 2$ term of the series expansion) to be

$$\Delta T_{\Delta s} \approx \Delta s^2 \frac{\sqrt{2m}}{4\sqrt{q}} E_s (s_0 E_s) \left[-1 + \frac{d}{S_0} \frac{1 - k_0^{-3/2}}{k_0 - 1} + \frac{3}{2} \frac{D}{s_0} k_0^{-5/2} \right], \quad (59)$$

$$\begin{aligned} \Delta T_{\Delta s} &\approx \frac{2.17nsec}{\sqrt{q}}, \\ \Delta T_\theta &\approx \frac{45nsec}{q}. \end{aligned} \quad (60)$$

For the smallest set of values for the voltages applied to the TOF grids, see values in Table 4, when $s = s_0$ and $k_0 \approx 1.577$; then we obtain that for $q = 1$ $\Delta T_{\Delta s} = 2.17nsec$ and $\Delta T_\theta = 45nsec$, and for $q = 7$ $\Delta T_{\Delta s} = 0.82nsec$ and $\Delta T_\theta = 6.43nsec$.

3 Data Analysis

3.1 Doubly Excited States of Helium Produced by Fast-Ion Impact

As we are interested in comparing our results with different theories, we first concentrate on the often investigated $2p^2(^1D)$ and $2s2p(^1P)$ states. The experimentally obtained energy spectra, integrated over all angles for different incoming ion charge states, were normalized to the same total number of ions hitting the Faraday cup. The peaks corresponding to the $2p^2(^1D)$ and $2s2p(^1P)$ states were least-square fitted to extract the Shore parameters [14] from the spectra as shown in Fig. 21.

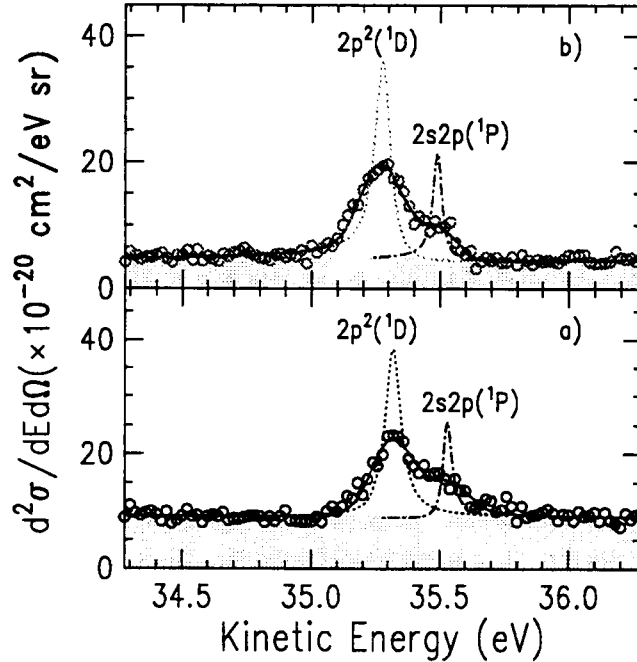


Figure 21: The least-square fit using a Fano profile of the electron spectra at 63° and 116° produced by C^{6+} ions. The dotted line is the profile of the $2p^2(^1D)$ resonance, the dashed-dotted line is that of the $2s2p(^1P)$ resonance and the solid line is the contribution of both lines convoluted with the resolution function of the spectrometer. The shaded area corresponds to direct ionization.

The theoretical line shape used in the fit procedure was [10]

$$\frac{d^2\sigma}{dEd\Omega} = C(E, \theta) + \sum_j \frac{A_j \epsilon_j + B_j}{\epsilon_j^2 + 1} \quad (61)$$

where E is the electron kinetic energy, θ is the electron emission angle with respect to the ion beam direction, $\epsilon = 2 \frac{E - E_{Rj}}{\Gamma_j}$, E_{Rj} is the energy of the j^{th} resonance, Γ_j is the natural line width, and A_j and B_j are the Shore parameters of the j^{th} resonance. $C(E, \theta) = D(\theta) + L(\theta) \times (E - E_{R1})$ is the contribution to the doubly differential cross section from the direct ionization process, $D(\theta)$ and $L(\theta)$ are arbitrary functions of θ . Throughout the fitting process, we assumed that the interaction of the emitted autoionized electron with the projectile does not significantly affect the measured profiles. Such post-collision interactions (PCI) effects are generally assumed to be important only when the projectile velocity is comparable to or smaller than the velocity of the emitted electron [6, 45]. We also assumed that state-to-state interference is very small compared to state-to-continuum interference and does not affect the analysis of the profiles.

The singly differential cross section for each resonance is obtained by integrating equation (61) over the electron energy

$$\frac{d\sigma}{d\Omega} = C(\theta) + \frac{\pi}{2} \Gamma_j B_j(\theta) \quad (62)$$

where $C(\theta) = \int C(E, \theta) dE$ is the contribution from the direct ionization. The line shapes and intensities of the peaks were obtained from the fit procedure using equation (61). The natural widths of the peaks were fixed in the fit with values of $0.072 eV$ for the $2p^2(^1D)$ state and $0.041 eV$ for the $2s2p(^1P)$ state [6, 46]. The fitting routine includes the spectrometer function that was earlier determined to be approximately a Gaussian profile with a full width half maximum of $0.15 eV$. The width used here was determined from the measured width of the Xe Auger lines $N_4 O_{23} O_{23}(^1S_0)$ and $N_5 O_{23} O_{23}(^1S_0)$ and correspond to accepted values. The energy spectrum of the Xe Auger lines $N_{45} O_{23} O_{23}$ with electron energies $27 \leq E_e \leq 37 eV$, shown in Fig. 22,

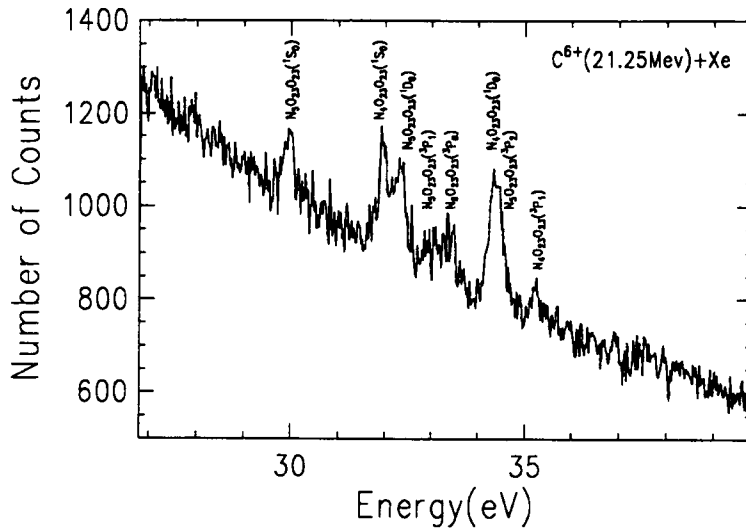


Figure 22: Xe spectra used for the energy calibration of the CMA.

was used to obtain the energy calibration for the CMA [47] given by the following equation:

$$\begin{aligned} V_{dac} &= \frac{98.6 - 66.1}{512} \times ChannelNumber + 66.1 \\ E(eV) &= 0.02546 \times ChannelNumber + 26.7859 \end{aligned} \quad (63)$$

where V_{dac} is the voltage set by the MP program that is applied to the inner and outer cylinder of the CMA through the programmable power supply and $ChannelNumber$ is the step number of a given energy scan. The images obtained on the PSD for the Xe Auger lines $N_4O_{23}O_{23}(^1S_0)$ and $N_5O_{23}O_{23}(^1S_0)$ which are known [47] to give an isotropic angular distribution, were used to obtain the angular detection efficiency of the microchannel plates (MCP).

From the Shore parameters we were able to obtain the singly differential cross section as a function of the emission angle using equation (62), which is shown in Fig. 23 for the $2p^2$ state. The θ dependence for the single differential cross section was obtained by fitting the data points shown in Fig. 23 with a function $F(\theta)$, where

$$F(\theta) = \alpha + \beta \cos \theta + \frac{1}{2} \gamma (3 \cos^2 \theta - 1). \quad (64)$$

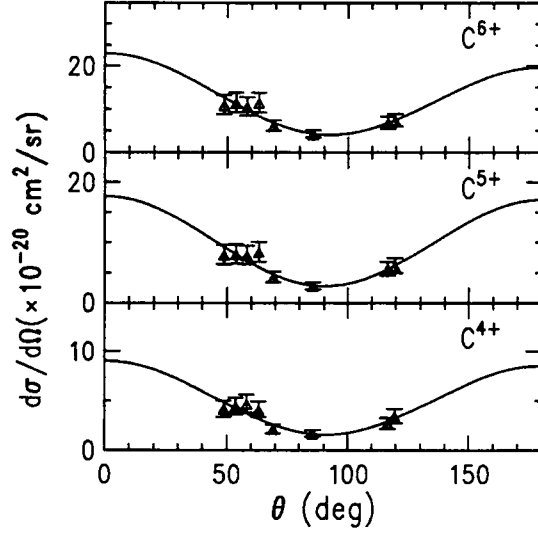


Figure 23: Singly differential cross section from the $2p^2(^1P)$ state as a function of the emission angle for different incoming C^{q+} ions. The solid line is the fitted curve to the data using the function $F(\theta)$.

The α , β and γ parameters obtained for the $2p^2(^1D)$ state are shown in Table 5.

A curve for the single differential cross section for angles between 0° to 180° was obtained, and shown in Fig. 23 with a solid line, for the $2p^2(^1D)$ state. The total cross section is then obtained by integrating $F(\theta)$ in the entire angular range as

$$\sigma_{tj} = \int F(\theta) d\Omega. \quad (65)$$

From the total cross section we are able to obtain its dependence on the incoming ion charge state that will give us some information about the mechanisms that lead to the production of the doubly excited states of He .

The singly differential cross section $d\sigma/d\Omega$ is obtained in absolute units by normalizing our result to the data from Giese *et al.* [6] at $\theta = 60^\circ$ for $\frac{d\sigma_{2p^2}}{d\Omega} + \frac{d\sigma_{2s2p}}{d\Omega}$. The main source of error in our results comes from the determination of the Shore parameters and the direct ionization contribution from the fit procedure of the $2p^2$ and

Table 5: Singly differential cross section fitting parameters for the $2p^2(^1D)$ state, in units of $10^{-20} \text{cm}^2/\text{sr}$.

C^{q+}	α	β	γ
C^{4+}	3.96	0.26	4.78
C^{5+}	7.66	0.29	9.72
C^{6+}	9.88	1.65	11.42

$2s2p$ states. As explained before, both the width and separation of the peaks were kept fixed during the fit. The widths used from Ref. [6, 46] have associated uncertainties leading to an $\sim 8\%$ error in the fitting parameters A_j , B_j , and $C(E, \theta)$. The separation between the peaks was obtained from the average separation obtained from the fit to all the spectra for different angles and incoming ion charge states. Then the data was fitted again keeping the separation between the peaks fixed to this average value. The uncertainty in the separation leads to $\sim 5\%$ uncertainty in the fitting parameters A_j , B_j , and $C(E, \theta)$. The statistical error for the measurements is less than 0.5%. The determination of the emission angle as an angular slice in the image obtained with the position sensitive detector has an uncertainty of $\pm 4^\circ$. All these errors are then propagated when the single differential and total cross sections are calculated.

3.2 Cascade-Decay Processes Produced in Photon-Atom Interactions

We measure the first step KLL Auger electrons emitted at different angles in coincidence with the subsequent LMM Auger electrons emitted from the decay cascade process in argon produced by photoexcitation by 3203.5eV photons. As explained in Section 2.2 the CMA was used to measure the KLL Auger electrons. A scan over the entire KLL group is shown in Fig. 24 where the KLL Auger lines are identified.

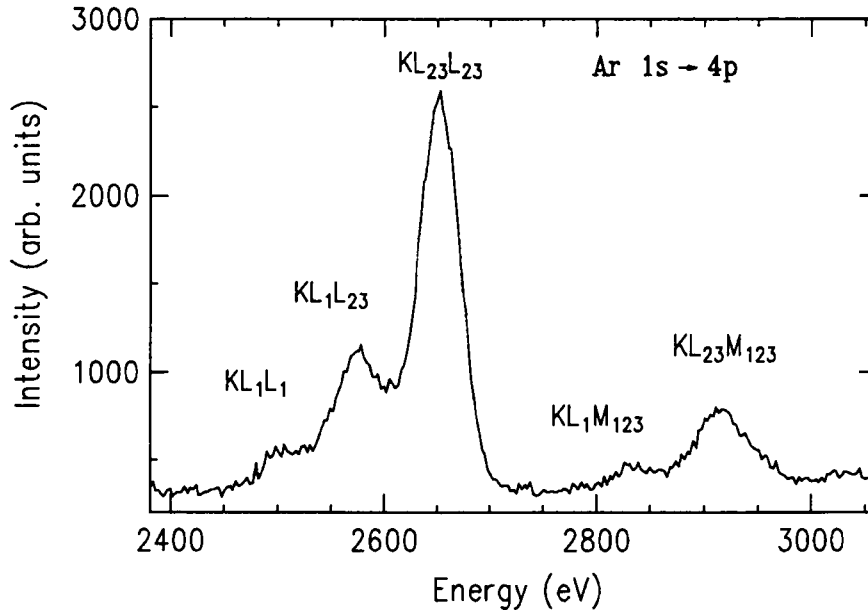


Figure 24: Spectrum of *KLL* Auger electrons for 3203.54 eV photon energy.

Subsequent *LMM* Auger electrons emitted in the decay cascade were measured in coincidence with the first step *KL₂₃L₂₃*, *KL₁L₂₃*, and *KL₂₃M₁₂₃* Auger electrons. The *LMM* Auger electrons were measured with the electron TOF spectrometer in the retarding mode with a retarding voltage of $-90V$. The different coincidence spectra are shown together with the non-coincidence spectrum for the entire *LMM* Auger group in Fig. 25. As can be observed in Fig. 25, the intensity of the different lines in the *LMM* Auger spectra change with the first-step *KLL* Auger electron selected for each coincidence spectrum. The energy resolution in this case is not sufficient to resolve the individual *LMM* lines. This is due to the fact that for even the strongest *KL₂₃L₂₃* Auger lines the coincidence rate was only about 22 mHz. To obtain reasonable statistics in the available time we had to make the measurements with low energy resolution.

A higher resolution *LMM* group is shown in Fig. 26. These spectra were obtained with the CMA working in its high resolution mode with pass energies of 100 eV and

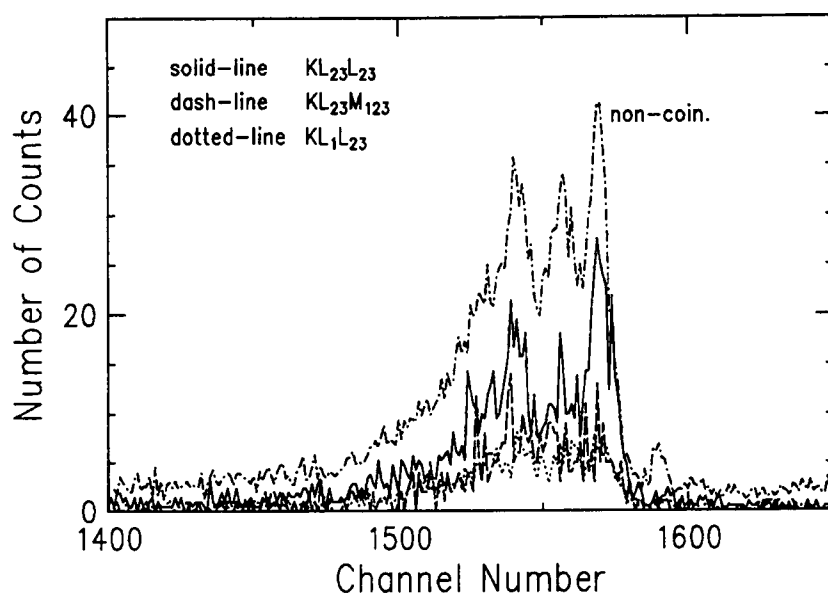


Figure 25: Time-of-flight spectrum of *Ar* LMM Auger electrons after *Ar* $1s \rightarrow 4p$ photoexcitation and a retarding voltage of $-90V$.

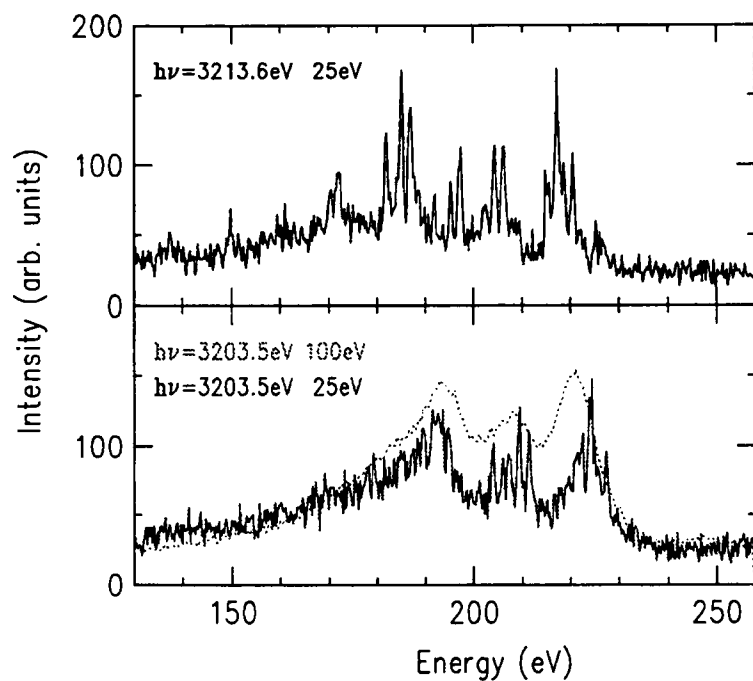


Figure 26: LMM Auger electron group. Top: 3213.6 eV; bottom: 3203.5 eV photon energy, for 25 eV and 100 eV CMA pass energies.

25eV for two different photon energies. It is possible to obtain a better energy resolution for the *LLM* group with both the CMA and TOF spectrometers. To obtain a better energy resolution with the TOF spectrometer a higher retarding voltage needs to be applied in the retarding stage to reduce the kinetic energy of the electrons. As can be seen in the TOF transmission curve (see Fig. 13 in Section 2.2) for low electron kinetic energies the transmission drops rapidly. Corrections to the electron spectra have been applied to take this effect into account.

As shown in Table 1, Section 1.2, in the second step of the decay process, low energy *LLM* Coster-Kronig transitions can occur together with *LLM* Auger transitions. To study the probability of this transition we looked at the emitted Coster-Kronig electrons, following photon excitation at the *Ar* $1s \rightarrow 4p$ resonance ($h\nu = 3203.5\text{eV}$). The Coster-Kronig electrons have a kinetic energy of around 40eV, an energy spectrum is shown in Fig. 27. The spectrum was taken with the CMA scanning a low electron energy region, at a photon energy of 3203.5eV.

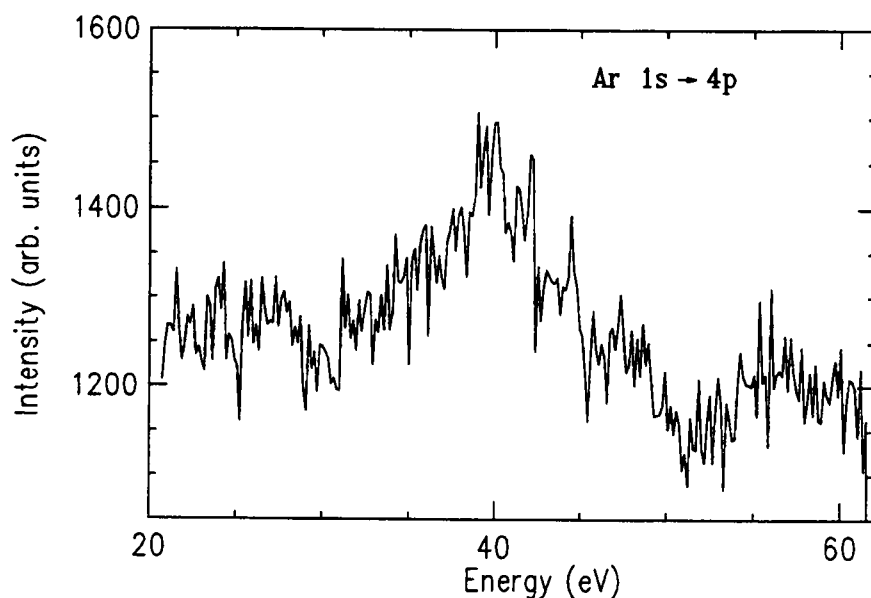


Figure 27: Argon Coster-Kronig electron spectra for 3203.54 eV photon energy.

If we want to compare the energy spectra obtained with the CMA and with the TOF spectrometer, we must first convert the TOF spectra that were measured as counts versus channel number into intensity versus energy in eV . The channel number in the TOF spectra represents the time of flight of the electron with respect to the reference signal that is given by the bunch marker signal of the synchrotron source. To convert the channel number into energy, it is not enough to just rescale the axis. We also need to convert the intensity of the signal from channel number to energy. We proceed as follows. First we fit the curve shown in Fig. 14 Section 2.2 with a function given by

$$y(x) = Prompt - \frac{1}{\sqrt{A + B(x + Acc)}} \quad (66)$$

where the y-value is the channel number, the x-value is the kinetic energy, Acc is the acceleration (or retarding) voltage, $Prompt$ is the position (in channel number) of the direct scattered light, A is the first parameter to describe the time to energy curve and B is the second parameter to describe the time to energy curve. This fitted curve is an approximation assuming

$$\begin{aligned} E_{kin} &= \frac{1}{2}mv^2 = \frac{m\Delta s^2}{2} \frac{1}{\Delta t^2} \\ \Delta t &\approx (p - ch) \\ E_{kin} &\approx \frac{1}{(p - ch)^2} \end{aligned}$$

where

ch = channel number,

E_{kin} = is the kinetic energy of the channel ch ,

p = is the channel number of the prompt.

This fitting procedure is not valid in the whole energy range, so we split the entire

energy range in two parts. The low energy part goes from 5 – 300eV and the high energy part is for energies > 80eV. For each of these regions we obtain a set of fitted parameters A , B , Acc and $Prompt$. Then with these parameters we can create the time to energy curve for the low and the high energy part. To obtain the right energy converted electron spectra we also need to obtain the right electron intensity as a function of energy. Then the intensity as a function of energy is obtained by doing the following conversion

$$I = \frac{z(ch) - \mu}{E(ch + 1) - E(ch)} \quad (67)$$

where

ch = channel number,

E = energy from the time to energy curve,

μ = background,

z = is the intensity in the channel ch of the raw data,

I = is the intensity for the energy $E(ch)$.

An electron spectrum taken with the electron TOF spectrometer after the time to energy conversion is shown in Fig. 28, where only the slower energy part of the spectrum is shown.

As we are interested to see if there is any angular correlation between the KLL and LMM Auger electrons we need some information about the efficiency of the MCP detector. As the KLL Auger electron emission is isotropic we used the angular distribution for the scan over the whole KLL group to obtain the efficiency of the MCP for each slice in angle taken in the spectrum. The efficiency factor is obtained for each slice in angle, and is then used as a correction factor to the energy spectra for each angular slice. The angular distribution for the KLL group is shown in Fig. 29. The angle in Fig. 29 is the angle ϕ measured on the MCP where 0° corresponds to the direction opposite to where the TOF spectrometer was mounted and 90°

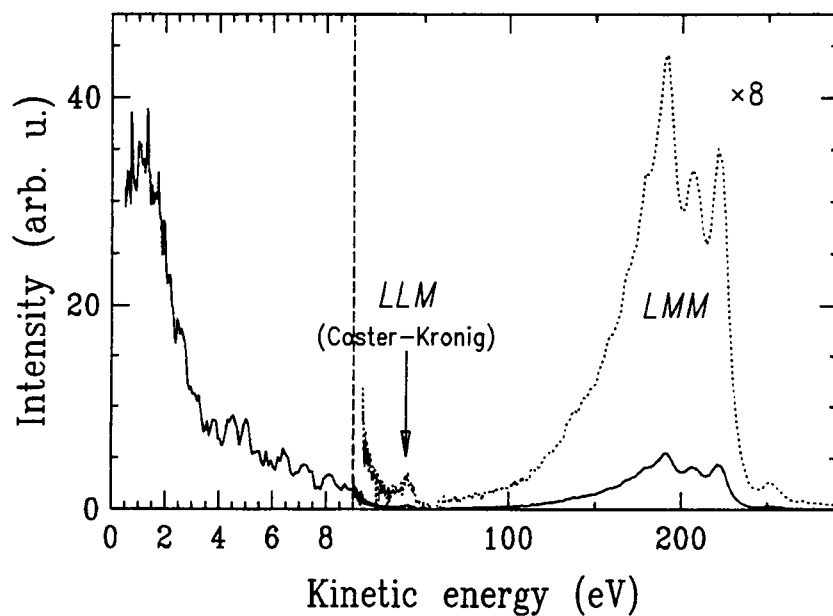


Figure 28: Electron spectrum taken with the electron TOF spectrometer after the time to energy conversion.

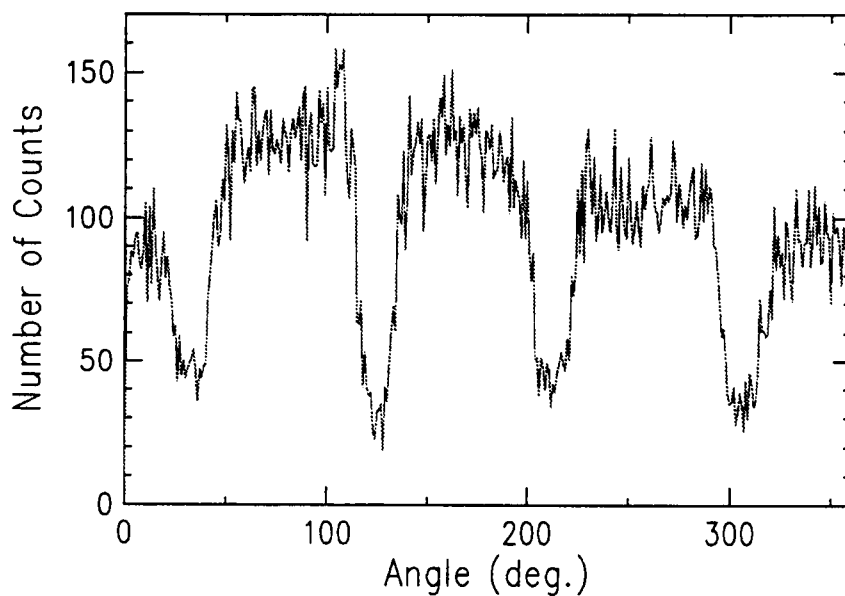


Figure 29: Angular distribution for the entire *KLL* Auger group.

corresponds to the photon beam direction. To obtain the angle θ of electron emission in the interaction region with respect to the TOF spectrometer direction (this will be the direction at which the *LMM* Auger electrons are detected) we need to use the following angle transformation

$$\theta = \cos^{-1}(-\cos \phi \sin 42.3^\circ) \quad (68)$$

Fig. 29 shows four regions where the signal is nearly uniform (or flat), illustrating the product of the isotropic angular distribution of the *KLL* Auger electrons with the detection efficiency of the MCP detector. The four minima in the spectrum occur at the location of the posts between the apertures in the inner cylinder of the CMA, as described in Section 2.1. These regions are shown as inactive on the PSD.

For different slices in angle we obtained the *LMM* Auger electron spectra to determine whether there is an angular correlation with the first-step *KLL* Auger electrons, for different groups of the *LMM* Auger lines. The determination of the emission angle has an uncertainty of $\pm 1^\circ$. As the angular distribution of the *LMM* group is symmetric with respect to the TOF spectrometer axis, half of the signal in the PSD detector is a replica of the other half.

The main source of error for the coincidence measurements is the statistical error due to the low count rate after taking different slices in angle and energy for the *LMM* Auger electron spectra. The statistical error was estimated to vary between 9% and 16% for the different regions considered in angle and energy. A second source of error comes from the determination of the angular efficiency factor due to the detection efficiency of the MCP detector. This error is also a statistical error and is estimated to be less than 4%.

3.3 Photoion Charge-State Distribution and Electron-Ion Coincidences

In establishing the dependence of the photoion charge state fractions on the different experimental parameters, we found that the charge state fractions depend on the voltage across the MCP detector, the pressure in the experimental chamber, and the thresholds of the constant fraction discriminator (CFD) shown in Fig. 18, Section 2.3. The ion time of flight spectrum for the different non-coincidence photoion charge states is shown in Fig. 30. The experimental parameters were set as follows: the CFD threshold was set at $400mV$, the pressure in the chamber was $2.5 \times 10^{-5}mbar$, the voltage on the upstream surface of the MCP detector was $-4350V$, and the voltage across the MCP detector was approximately equal to $1980V$.

As can be seen from Fig. 30 the peaks for the different charge states are very well resolved, so the fractions are obtained by calculating the area under each peak.

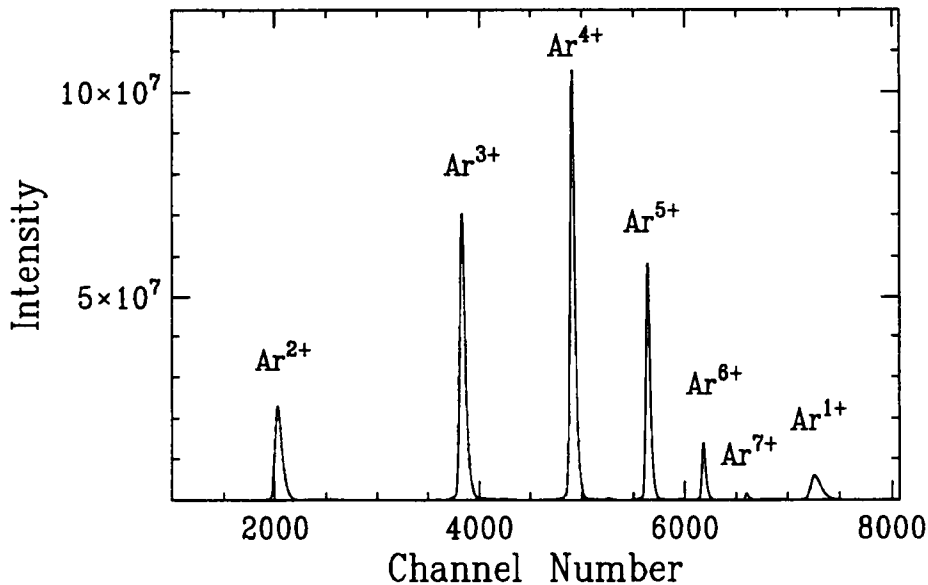


Figure 30: Charge state spectra obtained with the ion TOF.

The same is true for the coincidence spectra shown in Section 4.3.2. The photoion charge state fractions (or relative yields) for each charge state q are calculated from

$$\frac{I(Ar^{q+})}{\sum_{q'=1}^7 I(Ar^{q'+})} \quad (69)$$

where $I(Ar^{q+})$ is the intensity of the recoil Ar ion with charge state q . The two main sources of error in the determination of the charge-state fractions are the statistical error (bigger for the smaller peaks) and the determination of the starting and ending points in the calculation of the area for each peak (that will affect the background subtraction for each peak). The statistical error for all charge states is less than 3%. The error in the determination of the intensity of each photoion charge state, due to the method used to calculate the area under each peak, is less than 14% for all charge states.

4 Results

4.1 Doubly Excited States of Helium Produced by Fast-Ion Impact

From the data analysis described in Section 3.1 we were able to obtain the singly differential cross section for production of the the $2p^2(^1D)$ and the $2s2p(^1P)$ states for angles between 47° and 120° . Shown in Fig. 31 are the results for the $2p^2(^1D)$ state.

To compare our results with theory and previous data we need to examine the total cross section and its dependence on the incoming ion charge state. As discussed

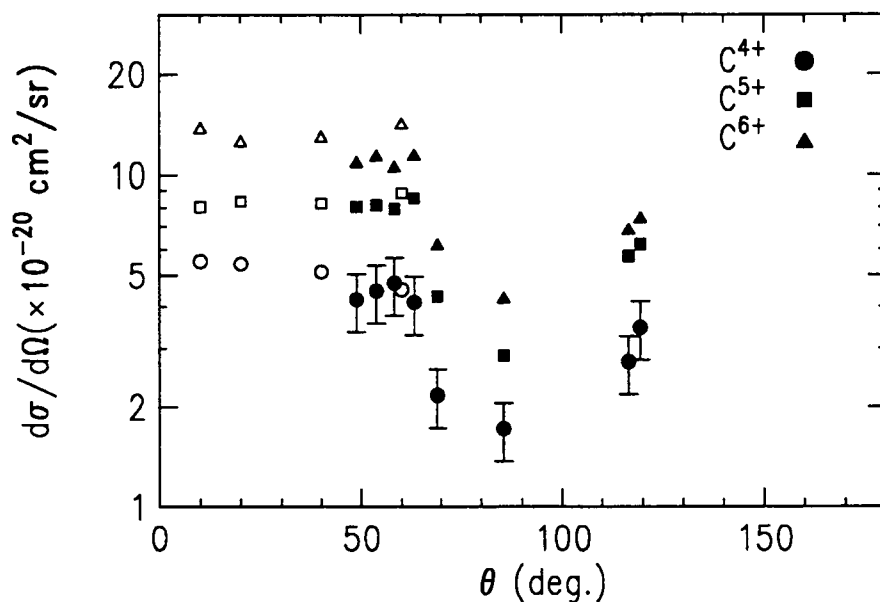


Figure 31: The differential emission cross section for autoionization from the $2p^2(^1P)$ state as a function of the electron emission angle after excitation by carbon ions. The solid symbols are the present data and the open symbols represent the data of Giese *et al.*[6]

in the Section 1.1, the total cross section ($\sigma_t(q)$) can be parametrized by

$$\sigma_t(q) = aq^2 + bq^3 + cq^4. \quad (70)$$

In a good approximation a simpler parametrization has been used to describe the q dependence of the total double-excitation cross section [6] by $\sigma_{tj} = Aq^B$.

As explained in Section 3.1 we were able to obtain the total cross section by integrating the singly differential cross section over all emission angles. The total cross section as a function of the incoming ion charge state is shown in Fig. 32 for the $2p^2(^1D)$ and the $2s2p(^1P)$ states.

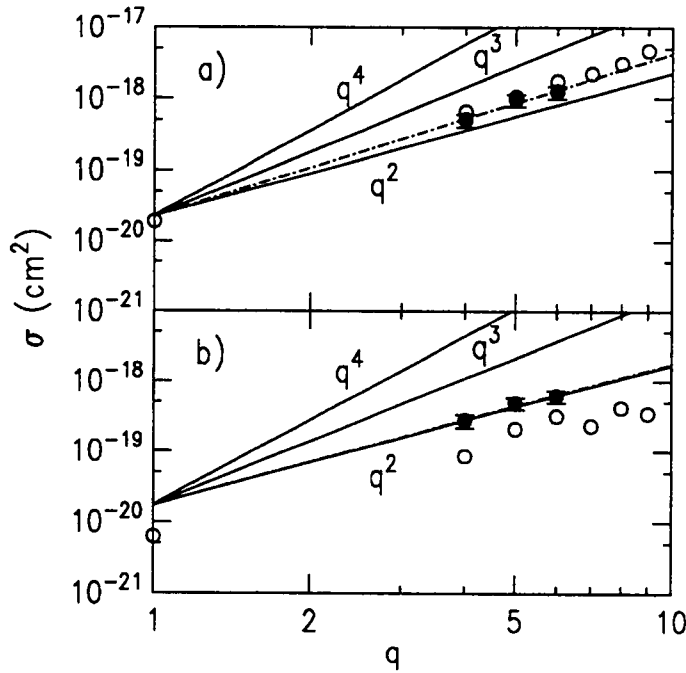


Figure 32: a) The total emission cross section from the $2p^2(^1D)$ state as a function of the projectile charge state; b) as in a), but from the $2s2p(^1P)$ state. The solid symbols are the present data; the open symbols are the data from Giese *et al.* [6]; the dashed line is a fitted curve to the data, using $\sigma_{tj} = A_j q^{B_j}$; the solid lines indicate the slopes corresponding to q^2 , q^3 and q^4 dependences.

The total cross section for each resonance as a function of the incoming ion charge state q is parametrized by $\sigma_{ij} = A_j q^{B_j}$. For the $2p^2(^1D)$ state the parameters are $A_{2p^2} = 2.25 \times 10^{-20} \text{ cm}^2$ and $B_{2p^2} = 2.27 \pm 0.18$. For the $2s2p(^1P)$ state $A_{2s2p} = 1.72 \times 10^{-20} \text{ cm}^2$ and $B_{2s2p} = 2.02 \pm 0.19$ (see Fig. 32). The data from Giese *et al.* [6] shown in this figure were obtained by assuming that the measured peak could be fitted by two theoretical peaks representing these two states, as Giese *et al.* did not resolve the peaks. From the present results we conclude that first-order processes are dominant for both states for $4 \leq q \leq 6$. These results are in good agreement with previous results [6] and calculations [10].

From the singly differential cross section we found that the production of the $2p^2(^1D)$ state cross section is approximately 1.85 ± 0.20 times more probable than that for the $2s2p(^1P)$ state for all angles $48.8^\circ \leq \theta \leq 119.5^\circ$ for C^{4+} ions and approximately 2.03 ± 0.22 times more probable for the C^{5+} and C^{6+} ions, as shown in Fig. 33. The ratio of the total double excitation cross sections for the production of the $2p^2(^1D)$ and the $2s2p(^1P)$ states is 1.97 ± 0.25 , averaged on $q = 4, 5, 6$. We were also able to independently determine the autoionization energies of the $2p^2(^1D)$ and the $2s2p(^1P)$ states to be $35.27 \text{ eV} \pm 0.04 \text{ eV}$ and $35.48 \text{ eV} \pm 0.04 \text{ eV}$, respectively. The energy difference between these states is $0.21 \pm 0.06 \text{ eV}$. These values are in good agreement with previous results from Refs. [58, 46] that give a value for the autoionization energies of 35.29 eV and 35.53 eV for these two states.

As can be seen from the discussion above we were able to compare the measured total cross section dependence on the incoming ion charge state with theoretical calculations. A more direct comparison of the measured electron emission cross sections with theory can be made if we look at the doubly differential cross section, and at the Shore parameters as a function of the electron emission angle for different incoming ion charge states. Then we can compare the measured doubly differential cross section for electron emission from the doubly excited states of helium with calculations done by Morishita *et al.* [11, 49]. In particular, we can examine the

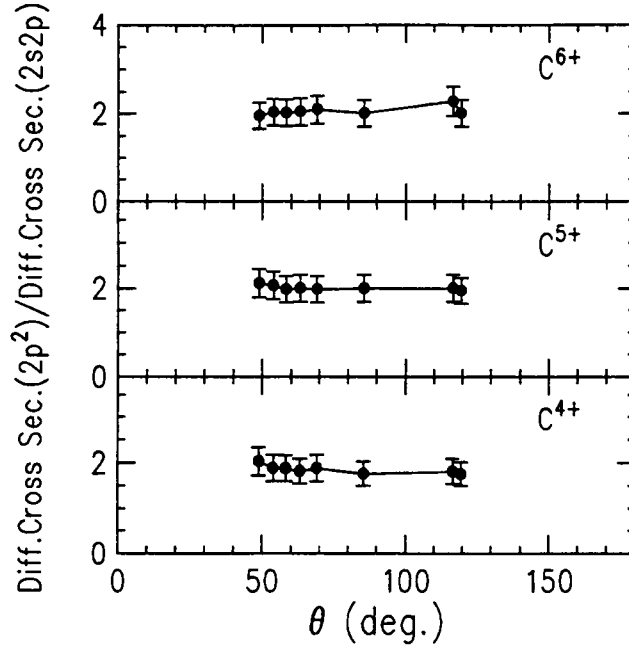


Figure 33: Ratio of singly differential cross sections for production of the $2p^2(^1D)$ state and the $2s2p(^1P)$ state, for different incoming ion charge states C^{q+} .

electron-emission spectra for electron kinetic energies from 32 eV to 40 eV. In these spectra we can identify and resolve the $2s^2(^1S)$, $2p^2(^1D)$, $2s2p(^1P)$ and $2p^2(^1S)$ states, as shown in Fig. 34, for 63° and 116° emission angles, for different incoming ion charge states. The peaks between 38 eV and 40 eV correspond to the contribution from the doubly excited states $2s3p(^1P)$, $2s3s(^1S)$, $2p3p(^1P)$, $2p3p(^1D)$, $2p3d(^1D)$, $2s4p(^1P)$, $2p3p(^1S)$, $2p3d(^1P)$, $2s4s(^1S)$ and $2p4p(^1D)$.

The data is compared to the theoretical calculations of Morishita *et al.* [11, 49] and Martín *et al.* [12] for C^{6+} ions for 63° and 116° emission angles, as shown in Fig. 35. From this figure we can see that theory and experiment agree quite well in the shape of the spectra, but not as well in the intensity of the resonances. The theoretical spectra [11, 49] are scaled to the experimental spectra so the direct ionization backgrounds coincide. The theoretical results need to be divided by 4.39 for the spectra at 63° and by 1.76 for the spectra at 116° . The differences between theory and experiment may be due in part to the way our data are normalized, and

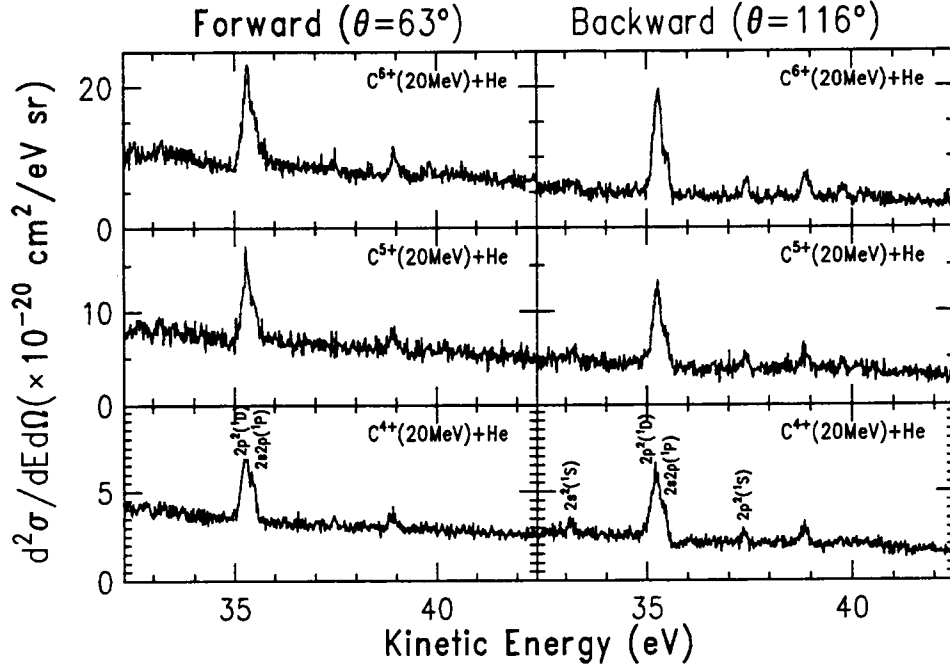


Figure 34: Doubly differential cross section for the production of electrons emitted at 63° and 116° with respect to the ion beam direction, with kinetic energies between 32 eV and 40 eV, for incoming carbon ion charge states $q = 4, 5, 6$.

in part to the fact that high angular momentum components L for the emitted electrons are not included in the calculation (maximum $L = 3$). We were also able to compare our data for C^{5+} and C^{4+} ions for emission angles of 63° and 116° with theoretical calculation done by Morishita *et al.* [49] as shown in Figs. 37 and 36. From the doubly differential cross section we can now obtain the Shore parameters for the $2p^2(^1D)$ and $2s2p(^1P)$ states for the experiment and the theoretical calculations done by Morishita *et al.* [49], and compare them with calculations done by Martín *et al.*[12]. We fitted the theoretical spectra with the same line shape we used for the experiment given by equation (61). We also obtained the Shore parameters for the theoretical spectra scaled to the experimental results, the parameters are plotted in Fig. 38 and listed in Table 6. The main cause of discrepancy in these results may be due to normalization differences between the calculated and the experimental values.

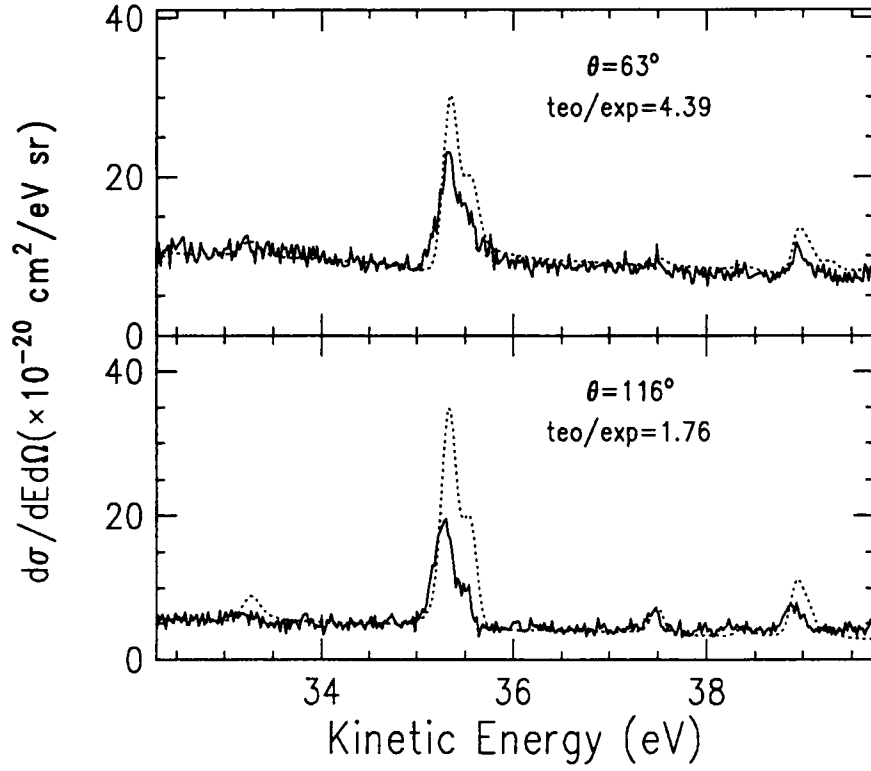


Figure 35: Experimental and calculated [49] doubly differential cross sections for the production of electrons emitted at 63° and 116° with respect to the ion beam direction, with kinetic energies between 32eV and 40eV for an incoming carbon ion charge state $q=6$. The solid line represents the experimental data and the dotted line the calculation.

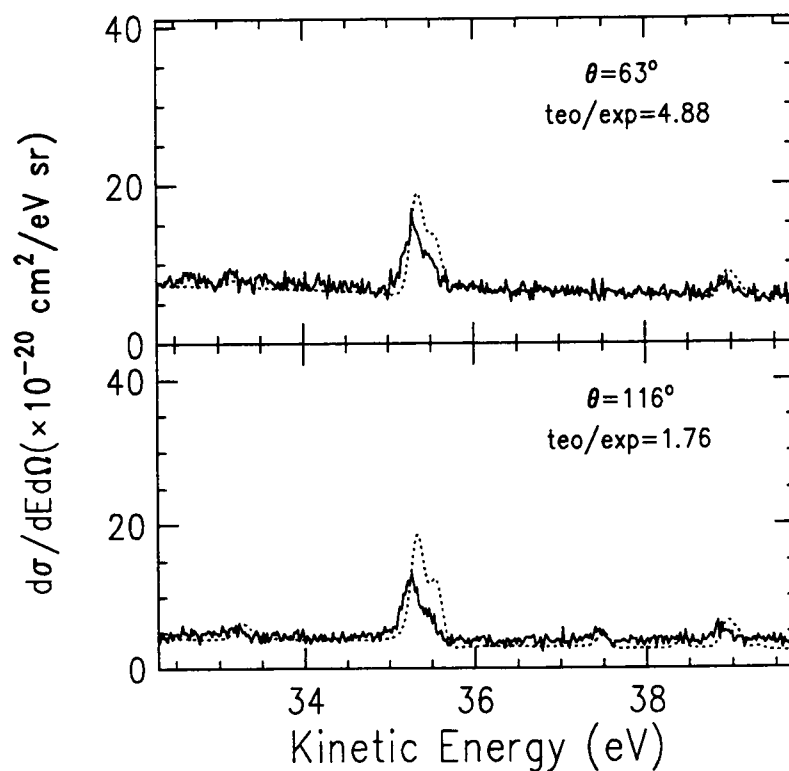


Figure 36: Experimental and calculated doubly differential cross sections for the production of electrons emitted at 63° and 116° with respect to the ion beam direction, with kinetic energies between 32eV and 40eV for an incoming carbon ion charge state $q=5$. The solid line represents the experimental data and the dotted line is the calculation [49].

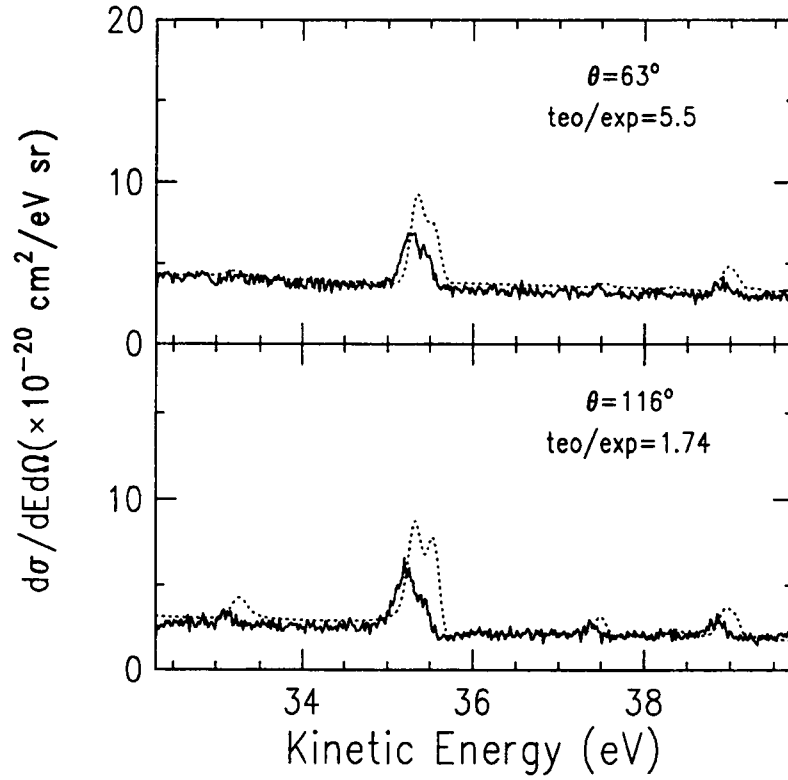


Figure 37: Experimental and calculated doubly differential cross sections for the production of electrons emitted at 63° and 116° with respect to the ion beam direction, with kinetic energies between 32eV and 40eV for an incoming carbon ion charge state $q=4$. The solid line represents the experimental data and the dotted line the calculation [49].

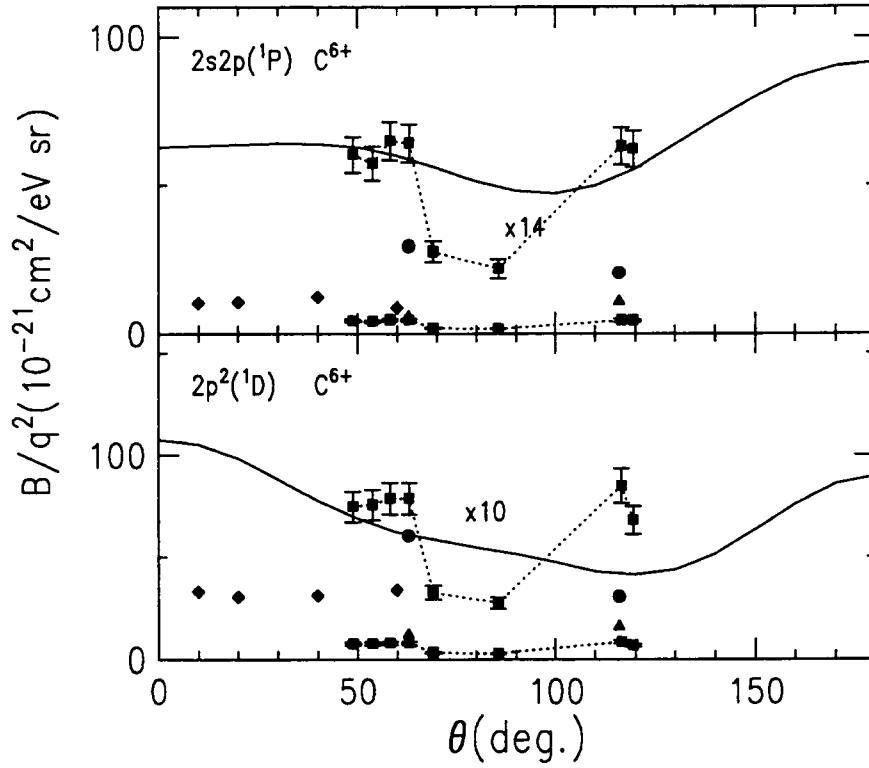


Figure 38: Experimental and calculated Shore parameters B/q^2 for electrons emitted between 0° and 180° with respect to the ion beam direction, for an incoming carbon ion charge state $q=6$ for the $2p^2(^1D)$ and $2s2p(^1P)$ states, respectively. The solid line shows the calculation from Martín *et al.* [12]. The solid squares represent the present experimental data, shown scaled with respect to the calculations of Martín *et al.* The solid circles give the results from the calculations of Morishita *et al.* [49], and the solid triangles represent the calculations scaled to the experimental data. The solid diamonds correspond to data from Giese *et al.* [6].

Table 6: The Shore parameter B/q^2 in units of $10^{-21} \text{cm}^2/\text{eV sr}$, for the production of the $2p^2(^1D)$ and $2s2p(^1P)$ states, for an incoming ion charge state $q = 6$.

Angle	Present Data	Giese <i>et al.</i> [6]	Calculated[49]
$2p^2(^1D)$			
10°		32.93	
20°		30.24	
40°		31.03	
49°	7.5 ± 0.9		
54°	7.6 ± 0.9		
58°	7.9 ± 0.9		
60°		33.72	
63°	7.9 ± 0.9		60.23
69°	3.3 ± 0.4		
86°	2.8 ± 0.3		
116°	8.5 ± 0.9		30.53
119°	6.8 ± 0.8		
$2s2p(^1P)$			
10°		10.22	
20°		10.48	
40°		12.29	
49°	4.3 ± 0.4		
54°	4.1 ± 0.4		
58°	4.6 ± 0.6		
60°		8.66	
63°	4.5 ± 0.5		29.69
69°	1.9 ± 0.3		
86°	1.6 ± 0.2		
116°	4.5 ± 0.5		20.61
119°	4.4 ± 0.5		

4.2 Cascade-Decay Processes Produced in Photon-Atom Interactions

In this section we discuss the angular correlation between $KL_{23}L_{23}$ Auger electrons emitted in the first step with the second step LMM Auger electrons obtained from the cascade Auger decay from Ar , initiated by photoexcitation to the $1s \rightarrow 4p$ resonance.

As was explained in Section 3.2, due to the low coincidence event rate we could only obtain sufficient statistics to study the coincidence between first step $KL_{23}L_{23}$ Auger electrons and the second step LMM Auger electrons. From the KLL Auger spectra we took different slices in angle and examined the coincident LMM Auger electron spectra, as shown in Fig. 39. In this figure the LMM Auger spectra are shown for $\theta = 48.9^\circ$, 80.3° , 99.9° and 130.5° .

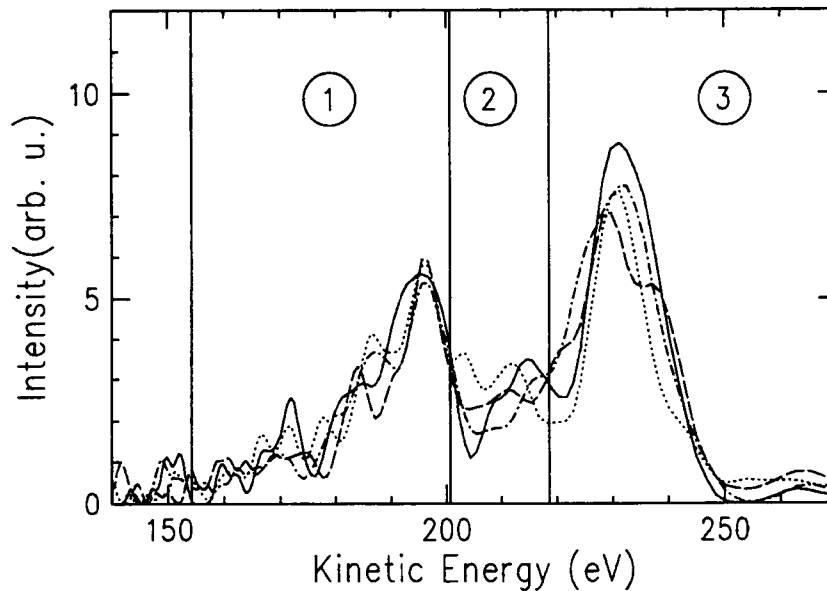


Figure 39: Energy spectra of LMM Auger electrons for different emission angles. The dashed line corresponds to $\theta = 48.9^\circ$; the dash-dotted line to $\theta = 80.3^\circ$; the dotted line to $\theta = 99.9^\circ$; and the solid line to $\theta = 130.5^\circ$. The vertical lines indicate the energy regions.

The three regions shown in Fig. 39 as (1), (2) and (3) correspond mainly to different LMM Auger lines. Previous experiments [51], where the $L_{23}MM$ Auger spectrum of argon was obtained by electron excitation, permitted assignments to some of the Auger lines. Comparing these three regions ((1), (2), and (3)) with the previous assignments we can see the following: for region (1) the $L_{23}M_1M_1$ are the more intense Auger lines; for region (2) the $L_{23}M_1M_{23}$ are the more intense Auger lines; and for region (3) the $L_{23}M_{23}M_{23}$ are the more intense lines. From Fig. 39 we can also see that the intensity of the LMM Auger lines in region (3) are very similar for all emission angles. For the three regions, shown in Fig. 39 as (1), (2) and (3), the intensity was obtained and plotted as a function of the $KL_{23}L_{23}$ Auger electron emission angle. The results are shown in Fig. 40, where three sets of data are plotted. Each set represent a different measurement of the same experiment where the $KL_{23}L_{23}$ Auger electrons were measured at different emission angles. This can be done by rotating the CMA around its symmetry axis. It can be seen from these results that regions (1) and (3) lie on a semicircle, meaning that their emission is isotropic when measured in coincidence with first step $KL_{23}L_{23}$ Auger electrons. However, region (2) shows a different angular dependence, showing evidence of correlation between the first step $KL_{23}L_{23}$ Auger electron and the second LMM Auger electron. This can be due to an alignment of the L vacancies corresponding to spectral region (2) after the $KL_{23}L_{23}$ Auger electron was emitted. The subsequent LMM Auger electron shows an angular distribution associated with this alignment.

The non-isotropic angular correlation observed for the distinct LMM slice (2) accompanies the detection of the first-step $KL_{23}L_{23}$ Auger electron in a particular direction, thereby breaking the symmetry so as to reveal the correlation. The KLL and the LMM Auger electron energies are very different. For this reason we believe that the large non-isotropic angular correlation observed represents an alignment effect (of the atomic system) rather than a dynamical post-collision interaction effect. As mentioned before, according to a previous assignment [51] of the LMM Auger

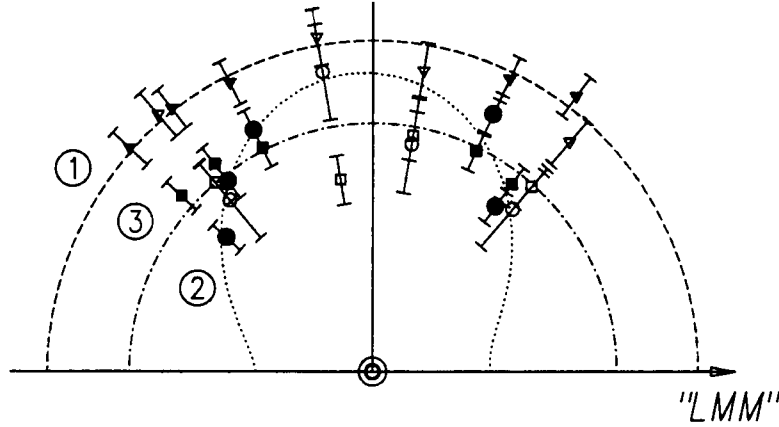


Figure 40: Intensity of LMM Auger electron emission in coincidence with $KL_{23}L_{23}$ Auger electrons, as a function of the $KL_{23}L_{23}$ Auger electron emission angle. The open and closed symbols represent different measurements. The triangles correspond to region (1), the circles to region (2), and the squares to region (3).

line, we can see that the lines contributing to slice (1) are mainly $L_{23}M_1M_1$ transitions; to slice (2) $L_{23}M_1M_{23}$ transitions; and to slice (3) $L_{23}M_{23}M_{23}$ transitions. In order to compare calculations with the experimental results, higher energy resolution must be achieved for the LMM Auger group. Due to the small value of the coincidence event rate for this experiment we were not able to run the TOF spectrometer at its highest energy resolution. A higher photon flux can compensate the transmission drop for lower electron kinetic energies, making the coincidence measurements possible at a better energy resolution and higher event rates. Such measurements are now becoming possible, using higher-flux beam lines developed more recently.

In Fig. 40 the dotted-line curve that best fit the angular distribution for slice (2) was obtained from the expression

$$f(\theta) = c \times (1 + \frac{\beta}{2}(3 \cos^2(\theta - \lambda) - 1)) \quad (71)$$

where the degree of linear polarization is ~ 1 , $c = 141.47$ is the amplitude in arbitrary units, $\lambda = 3^\circ$ is the tilt of the polarization axis in degrees (for measurements in the plane perpendicular to the photon beam) and $\beta = -0.5$ is the angular distribution parameter β . Where the numerical values given are the results of the least-square fit to expression (71). As discussed in Section 1.2.2, this expression for the angular distribution was obtained for the emission of the photoelectrons in the dipole approximation. Here it is used to fit the angular distribution for an emission of a particular group of the *LMM* Auger electrons, measured in coincidence with the first step $KL_{23}L_{23}$ Auger electrons.

4.3 Photoion Charge-State Distribution and Electron-Ion Coincidences

4.3.1 Photoion Charge-State Distribution

In these experiments, we measured ion-charge-state distributions as a function of the photon energy for an argon gas target. The values of the measured charge-state fractions depend on several parameters such as the voltage across the MCP, the pressure in the experimental chamber, and the thresholds of the constant-fraction discriminator (CFD) as discussed in Section 2.3. We concentrated our studies on these dependences, to assure accurate values of the measured charge-state fractions.

First we examine the dependence on the CFD thresholds. In this case we set the photon beam energy to 3203.5 eV (*Ar* $1s \rightarrow 4p$ resonance), the pressure of the experimental chamber to $2.5 \times 10^{-5} \text{ mbar}$, the voltage in the first surface plate of the MCP to $-4350V$, with the potential difference across both MCPs set equal to 1980V. The results are shown in Fig. 41. From Fig. 41 we can estimate the variation of the different charge state fractions due to a change in the threshold value of the

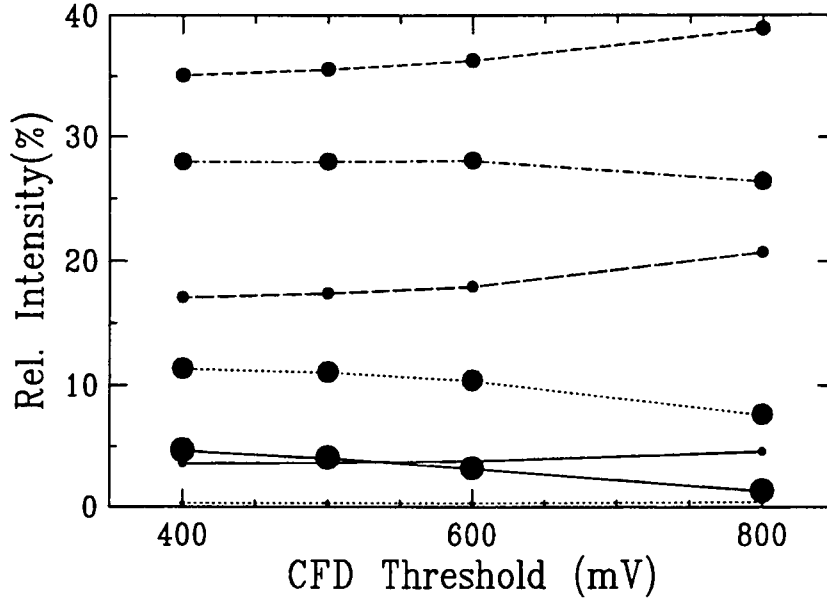


Figure 41: Ion charge-state fractions as a function of the CFD threshold. The largest symbol corresponds to Ar^{1+} , and the symbols decrease in size with increasing charge state, with the smallest symbol corresponding to Ar^{7+} .

CFD, as shown in Table 7. We can see that the biggest effect, for $600mV$, is suppression of the smallest charge-state fractions $q = 1, 2$. At $600mV$ the charge-state fraction for $q = 1$ changes more than 30% with respect to the value at $400mV$. This is due to the fact that the pulse height from the MCP detector is smaller for these charge states than for $q \geq 3$. If we continue increasing the CFD threshold, to $800mV$, bigger pulses are discriminated leading to a suppression of charge-state fraction $q = 1, 2, 3$ and enhancing the fractions for $q = 4, 5, 6, 7$. The values of the ion-charge-state fractions at $500mV$ and $400mV$ are measured to be the same within the experimental error. If we set the threshold voltage of the CFD at a value smaller than $500mV$ the relative ion-charge-state fractions do not change.

Next we fixed the CFD threshold to $600mV$, kept the same photon energy, the same pressure in the experimental chamber, and measured the charge-state fractions as a function of the voltage across the MCP detector, and the voltage applied in the entrance surface of the MCP detector. The results are shown in Fig. 42. As can be

Table 7: Charge-state fractions for different CFD threshold voltages with the relative differences in the fractions at 600mV and 800mV given with respect to the values at 400mV (with $h\nu = 3203.5\text{eV}$ and MCP front plate at -4350V).

Ar^{q+}	400mV	600mV		800mV	
		Rel. Frac. %	Rel. Dif. %	Rel. Frac. %	Rel. Dif. %
Ar^{1+}	4.66 ± 0.3	3.18	31.8	1.31	71.9
Ar^{2+}	11.32 ± 1.0	10.38	8.3	7.60	32.9
Ar^{3+}	27.93 ± 1.2	28.04	0.4	26.40	5.5
Ar^{4+}	35.08 ± 1.1	36.31	3.5	38.92	10.9
Ar^{5+}	17.08 ± 1.3	17.94	5.0	20.72	21.3
Ar^{6+}	3.61 ± 0.2	3.83	6.1	4.62	27.9
Ar^{7+}	0.32 ± 0.1	0.33	3.1	0.44	37.5

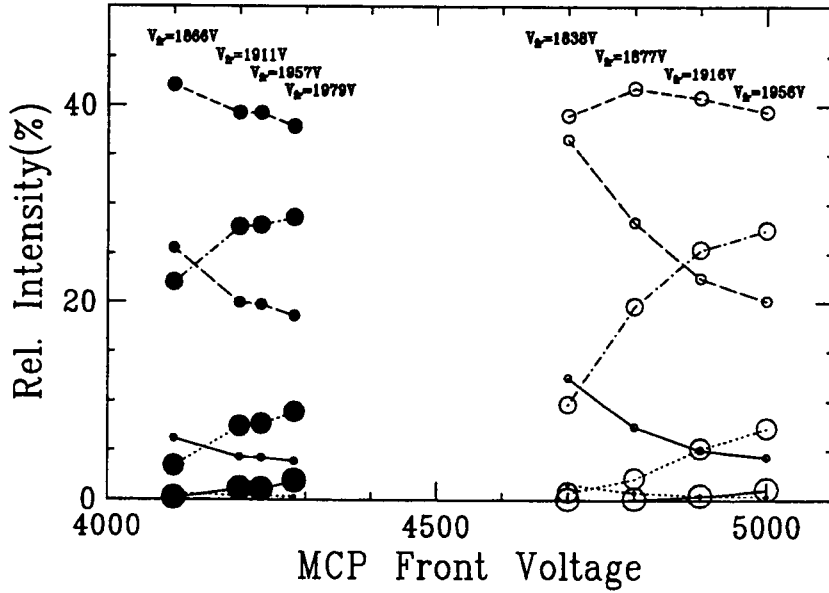


Figure 42: Ion charge-state fractions as a function of the MCP front voltage for different potential differences across the MCP detector. V_f denotes the potential difference applied across both MCPs. The largest symbols correspond to Ar^{1+} and decrease in size with increasing charge state.

seen in this figure the measured ion-charge-state fractions mostly depend on the voltage applied across the MCP detector and are approximately independent of the voltage applied to the entrance surface of the MCP detector. It is also seen that we were not able to reach to a region where the ion-charge-state fully saturates with the voltage applied across the MCP. The higher voltages needed to achieve saturation cannot be reached because of the 2000V restriction specified by the manufacturer.

To correct our measured values of the charge-state fractions due to this partial, but incomplete saturation effect we fitted the curves in Fig. 42 by an exponential function given by

$$Fraction = ampl \times (1 - \exp((Voltage - offset)/const)). \quad (72)$$

For a photon energy of $h\nu = 3203.5eV$ the measured and corrected values for the ion-charge-state fractions are shown in Table 8, from which we can see that the difference between the measured and the corrected fractions for the experimentally set and extrapolated voltages across the MCP detector is less than 9% for all the

Table 8: Measured and corrected values of the charge-state fractions for an applied potential difference across the MCP detector of $V_{fr} = 1979V$, corrected to $V_{fr} = 2730V$ (with $h\nu = 3203.5eV$, CFD Threshold Voltage = 400mV).

	Measured Values	Corrected Values
	Rel. Int. %	Rel. Int. %
Ar^{1+}	3.02 ± 0.3	2.78
Ar^{2+}	11.09 ± 1.0	11.56
Ar^{3+}	27.93 ± 2.1	28.72
Ar^{4+}	35.43 ± 2.7	33.46
Ar^{5+}	17.35 ± 1.3	17.01
Ar^{6+}	3.68 ± 0.3	3.78

charge state fractions. Thus the main source of error in the determination of the values of the charge-state fractions remains the fact that the maximum voltage difference that can be applied across the MCP detector is 2000V.

Finally, we studied the ion-charge-state fraction dependence on the pressure inside the experimental chamber. The pressure in the chamber was varied between $7.5 \times 10^{-7} \text{ mbar}$ to $3.4 \times 10^{-5} \text{ mbar}$. In this case the CFD threshold was 400mV, the energy of the photon beam was 3203.5eV, as before, and the voltages applied to the entrance surface of the MCP detector and across the MCP detector were -4350V and 1979V, respectively. The results, shown in Fig. 43, prove that the ion-charge-state fractions do not depend significantly on the pressure in the experimental chamber. After completing these studies of the ion-charge-state fraction dependence on all

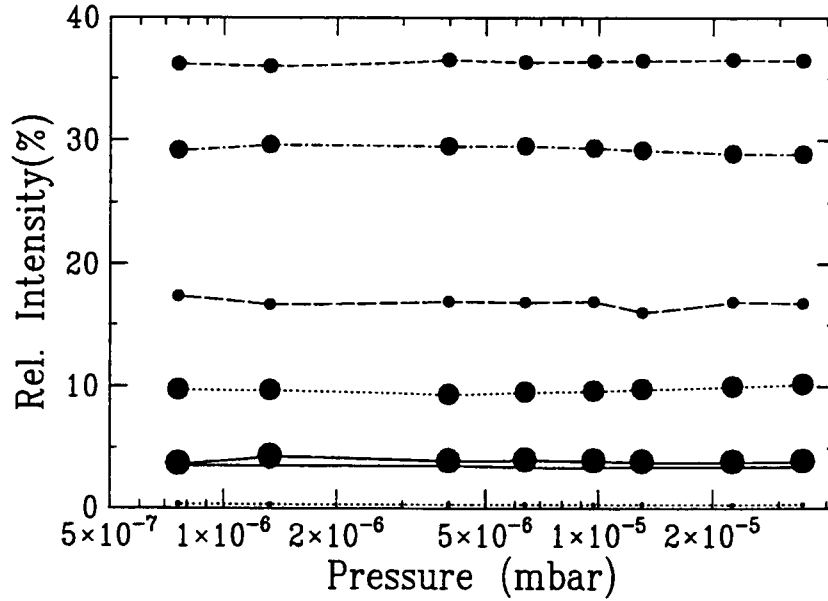


Figure 43: Ion charge-state fractions as a function of the pressure in the experimental chamber. The largest symbols correspond to Ar^{1+} , and decrease in size with increasing charge state with the smallest symbol for Ar^{7+} .

the significant variable experimental parameters we could proceed to measure the fractions as a function of the photon beam energy. We chose the following final setting for the experimental parameters. The CFD threshold was set at $400mV$; the pressure in the chamber was $2.5 \times 10^{-5}mbar$; the voltage applied to the entrance surface of the MCP detector was $-4350V$; and the potential difference across the MCP detector was approximately $1980V$. The time-of-flight spectrum obtained with the ion TOF spectrometer for a photon energy of $3203.5eV$ is shown in Fig. 44, where the ion charge states are identified. We thus obtained the ion-charge-state fractions from the time-of-flight spectra as a function of the photon energy $3190eV \leq h\nu \leq 3220eV$, with the results shown in Fig. 45. The measured fractions for Ar^{7+} are less than 1% for all photon energies.

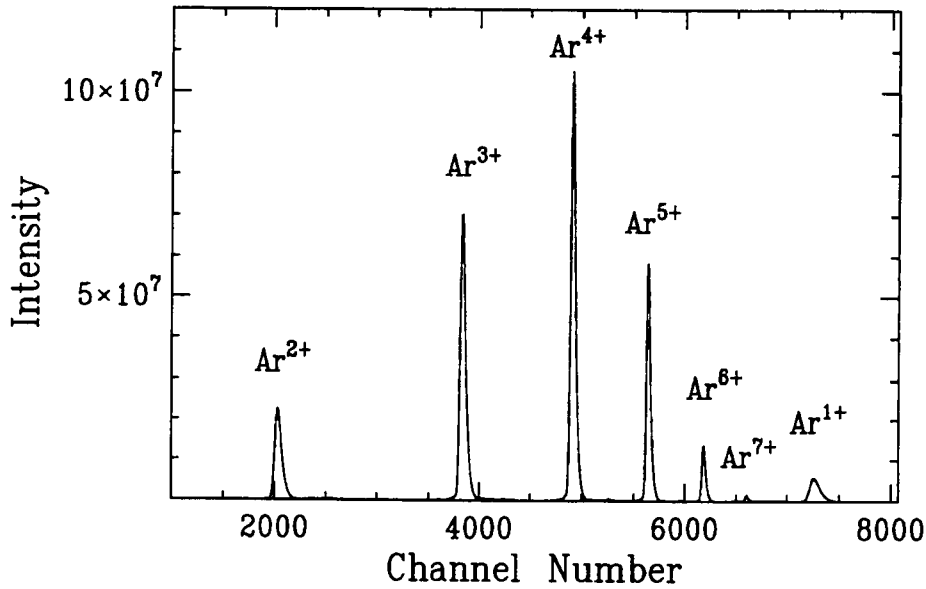


Figure 44: Charge-state spectra obtained with the ion TOF.

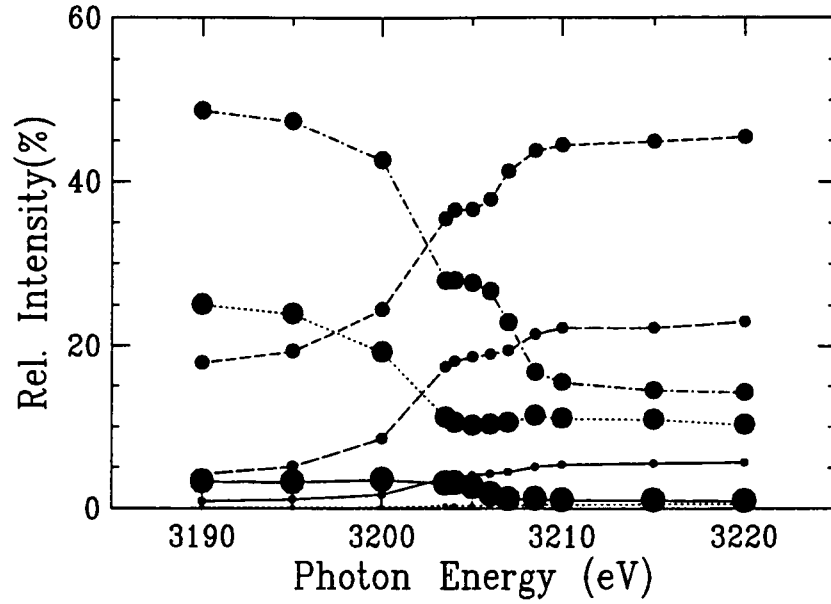


Figure 45: Ion charge-state fractions as a function of the photon-beam energy. The largest symbols correspond to Ar^{1+} , and decrease in size with increasing charge state, with the smallest symbol corresponding to Ar^{7+} .

If we look at the ion-charge fractions at energies above the K -threshold (3206.3eV), for example at more than above 3210eV photon energy, we compare our results with previous measurements done by Ueda *et al.* [18] and Doppelfeld *et al.* [31]. The values are shown in Table 9, where it is seen that the biggest differences between the different measurements are for the lower charge states. As we have seen if the potential difference across the MCP detector is not large enough the measured charge state fractions for $q = 1, 2$ are measured to be smaller than the true ones. Consequently, for charge-state fractions $q = 3, 4, 5, 6, 7$ the measured values will be systematically too large. A similar effect is obtained if the signal from the detector is passed through a CFD whose threshold value is set at a high value. In this case the measured charge-state fractions for $q = 1, 2, 3$ will be artificially smaller and for $q = 4, 5, 6, 7$ artificially larger than the true values. In Table 7 we estimate the differences in the charge-state fractions measured if the threshold of CFD was set at 800mV and not at 400mV .

Table 9: Asymptotic values of the relative intensities of Ar^{q+} ion spectra for photon energies above 3210eV, ^a Ueda *et al.*(1991) [18], ^b Doppelfeld *et al.*(1993) [31].

	Present Work	Other Work	
	Rel. Int. %	Rel. Int. % ^a	Rel. Int. % ^b
Ar^{1+}	0.96 ± 0.1	1.8	0.58
Ar^{2+}	10.54 ± 0.9	11.7	8.50
Ar^{3+}	14.34 ± 1.3	10.9	11.52
Ar^{4+}	45.31 ± 4.1	40.4	44.85
Ar^{5+}	22.73 ± 2.0	24.7	26.97
Ar^{6+}	5.56 ± 0.5	10.5	7.58
Ar^{7+}	0.56 ± 0.1	< 1%	< 1%

If we look at the ion charge fractions at energies above the K -threshold (3206.3eV) we see from Table 1 that Ar^{1+} and Ar^{2+} cannot be produced through a K -shell Auger processes (although the KMM Auger transition produces Ar^{2+} , its contribution is less than 1% of the total ion intensity). If a radiative decay of the K -hole occurs with the emission of K_{α} X-rays, a new hole is created in the L_{23} subshell and the subsequent $L_{23}MM$ Auger transition would produce Ar^{2+} [18]. If K_{β} is emitted instead of K_{α} , the resultant ion would be Ar^{1+} , the possible processes that lead to Ar^{1+} and Ar^{2+} are shown in Table 10.

The sum of the relative intensities of Ar^{1+} and Ar^{2+} is therefore expected to give the fluorescence yield for the Ar K -hole, if we neglect the shake-off and the KMM Auger transitions. The shake-off would suppress the Ar^{2+} production by about 10% [53] and the KMM Auger transitions would enhance the Ar^{2+} production by about 10% [52], so these contributions are presumed to cancel out. The asymptotic values of the relative intensities of Ar^{q+} for $q = 1...7$ above 3210eV are shown in Table 9. From the asymptotic values of the relative intensities of Ar^{1+} and Ar^{2+} we

Table 10: Radiative decay processes for a K -hole produced by the photoionization of Ar, leading to Ar^{1+} and Ar^{2+} for photon energies above the Ar K edge.

Initial Hole	Subsequent processes	Subsequent processes	Produced ions photoionization
[K]	[M] + ω_1		Ar^{1+}
	[L] + ω_2	[M] + ω_3	Ar^{1+}
[K]	[L] + ω_4	[MM] + e^-	Ar^{2+}
		[MMM] + $2e^-$	Ar^{3+}
	[MM] + e^-		Ar^{2+}

have estimated the fluorescence yield to be 0.115 ± 0.02 , and the Auger yield to be 0.885 ± 0.02 . The value for the fluorescence yield is in reasonable agreement with the value of 0.135 estimated by Ueda *et al.* [18] and with the compiled value of 0.118 by Krause *et al.* [54].

From Fig. 45 we obtained the ion-charge-state fractions as a function of the photon energy, normalized to the total ion yield measured by Armen *et al.* [21], shown in Figs. 46 and 47. From the normalized total ion-charge-state fractions alone it is difficult to extract detailed information about the decay path ways due to the complicated decay processes that are taking place. Coincidence measurements of the ion charge state fractions with individual first step KLL Auger electrons selected can help to understand the cascade-decay processes as described by Armen *et al.* [21]. We were able to make coincidence measurements having the $KL_{23}L_{23}$ and KL_1L_{23} Auger electrons as the first step Auger decay at two photon energies $h\nu = 3203.5\text{eV}$ and 3213.5eV ; the results are shown in Section 4.3.2.

The average charge state as a function of the photon energy can be obtained from the ion-charge-state fractions as

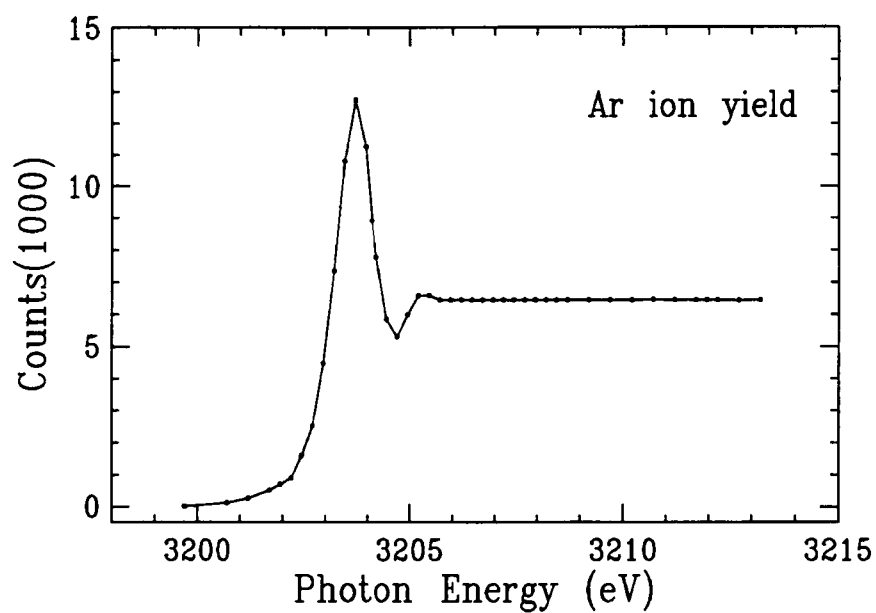


Figure 46: Measured total ion yield [21].

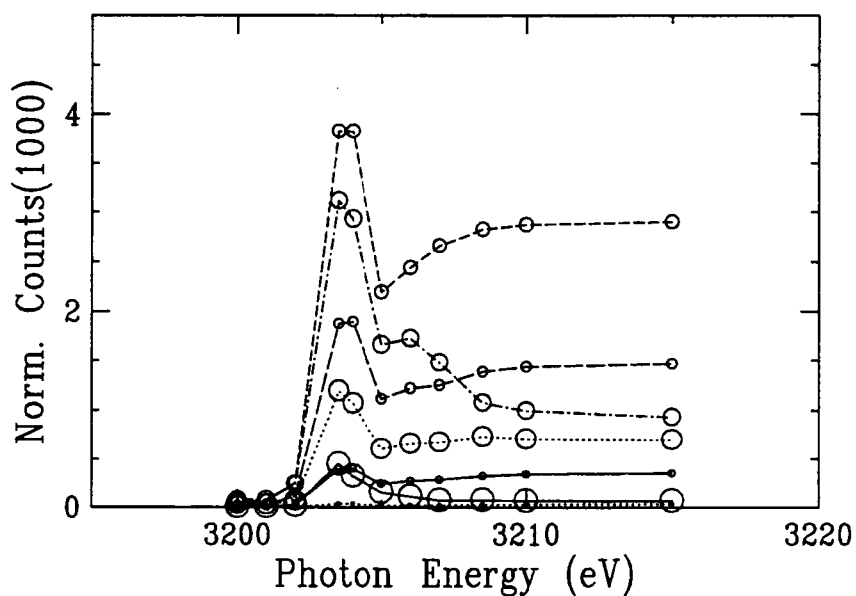


Figure 47: Relative ion-charge-state yields as a function of the photon-beam energy normalized to the total ion yield. The largest symbol corresponds to Ar^{1+} , and decrease in size with increasing charge state, with the smallest symbol for Ar^{7+} .

$$\bar{q} = \frac{1}{100} (q_1 P_1 + q_2 P_2 + q_3 P_3 + q_4 P_4 + q_5 P_5 + q_6 P_6 + q_7 P_7) \quad (73)$$

where $q_i = 1 \dots 7$ and P_i = percentage of each charge state q_i and is shown in Fig. 48. From this figure we can see that beyond the *Ar* *K*-shell threshold, the spectrum still keeps changing until 6eV above threshold. This transition behavior above threshold is due to recapture of slow photoelectrons by post-collision interaction (PCI) with a *KLL* Auger electron [17, 18, 31]. For photon energies below the *K* threshold the photoelectron is excited to a discrete (though broad) *np* state. As a spectator to Auger and radiative transitions, it may be ejected by shake-off during the cascade. The smooth change observed in all ionic charge state abundances, and observed as well in the $\bar{q}$ from the $1s \rightarrow 4p$ resonance up to 6eV above threshold, is caused by a balance between loss of the *np* electron on one hand, and recapture of the photoelectron on the other.

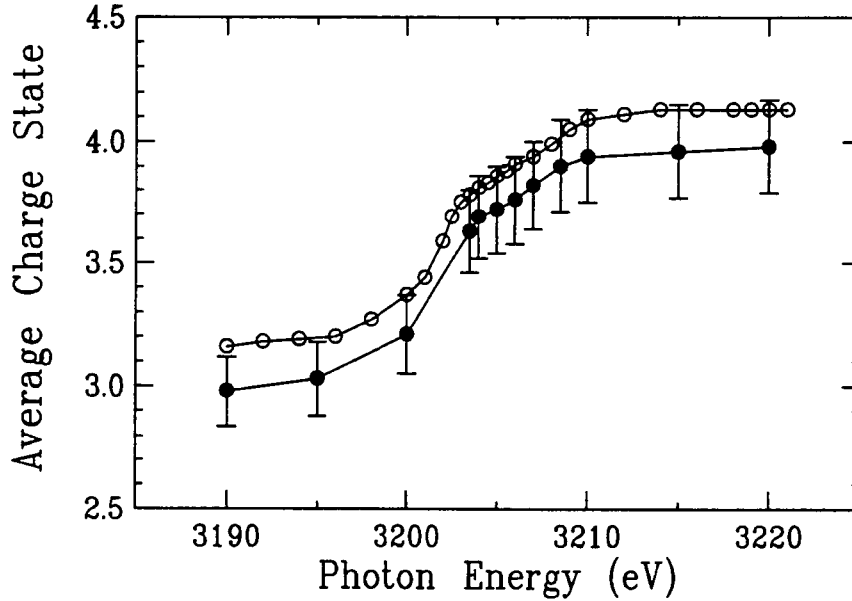


Figure 48: Average charge state as a function of the photon-beam energy. The solid circles correspond to present data and the open circles represent the measurements from Doppelfeld *et al.* [31].

4.3.2 Electron-Ion Coincidences

In these experiments we measured the ion-charge-state distributions in coincidence with angle resolved $KL_{23}L_{23}$, KL_1L_{23} Auger electrons at 3203.5eV (Ar $1s \rightarrow 4p$ resonance) and 3213.5eV photon energies for an argon gas target. Our principal interest lies in measuring the ion-charge-state fractions for different KLL Auger electrons as the first step in an Auger decay process, to obtain information about which decay paths dominate for a particular first step KLL Auger electron. The decay path, and consequently the charge-state fractions, depends on the photon energy and on the different possible decay paths, as shown in Table 1, Section 1.2. The time-of-flight spectra for ions in coincidence with different KLL Auger electrons are shown in Figs. 49, 50, 51, 52. The relative ion-charge-state fractions for the different coincidence spectra are shown in Fig. 53.

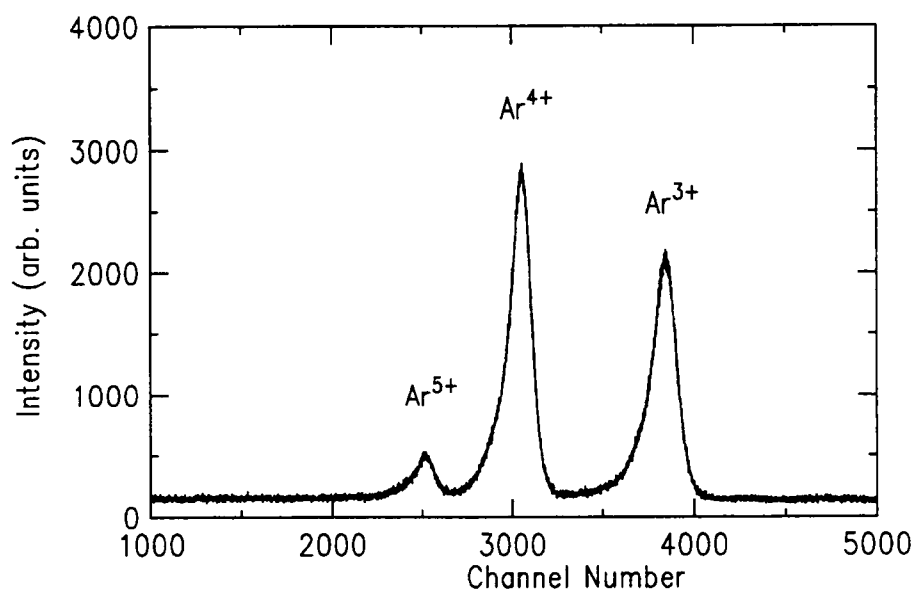


Figure 49: Ion time-of-flight spectrum observed in coincidence with $KL_{23}L_{23}$ Auger electrons, for 3203.5 eV photon energy.

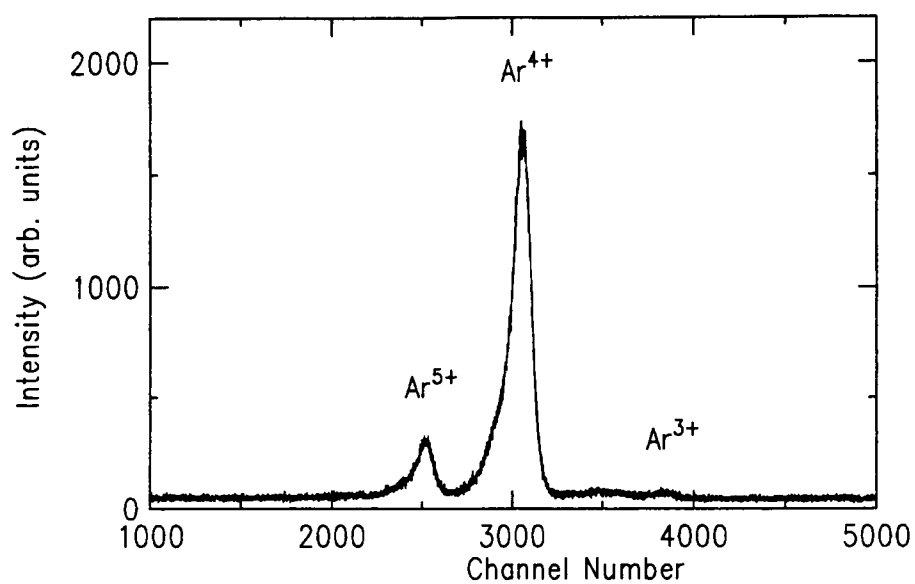


Figure 50: Ion time-of-flight spectrum observed in coincidence with $KL_{23}L_{23}$ Auger electrons, for 3213.5 eV photon energy.

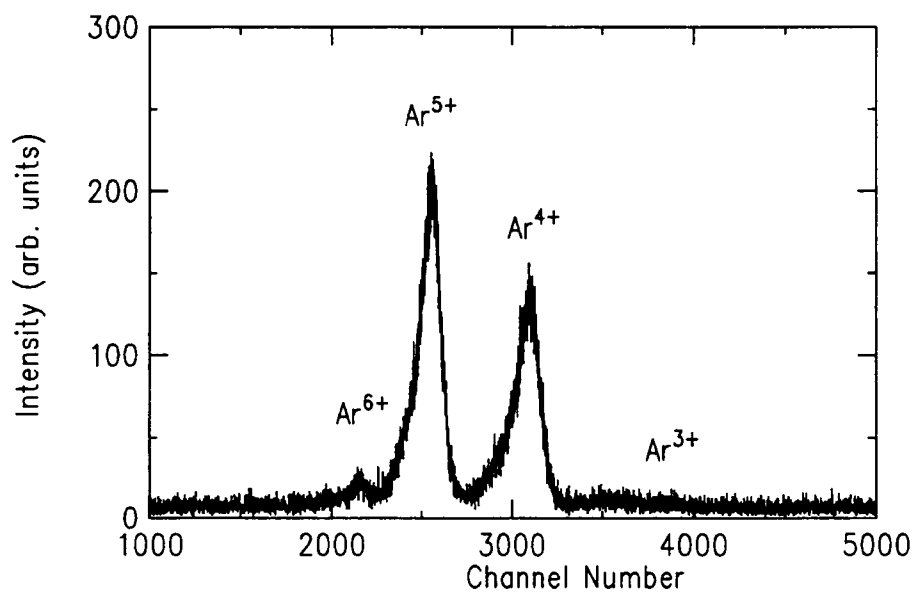


Figure 51: Ion time-of-flight spectrum observed in coincidence with KL_1L_{23} Auger electrons, for 3203.5 eV photon energy.

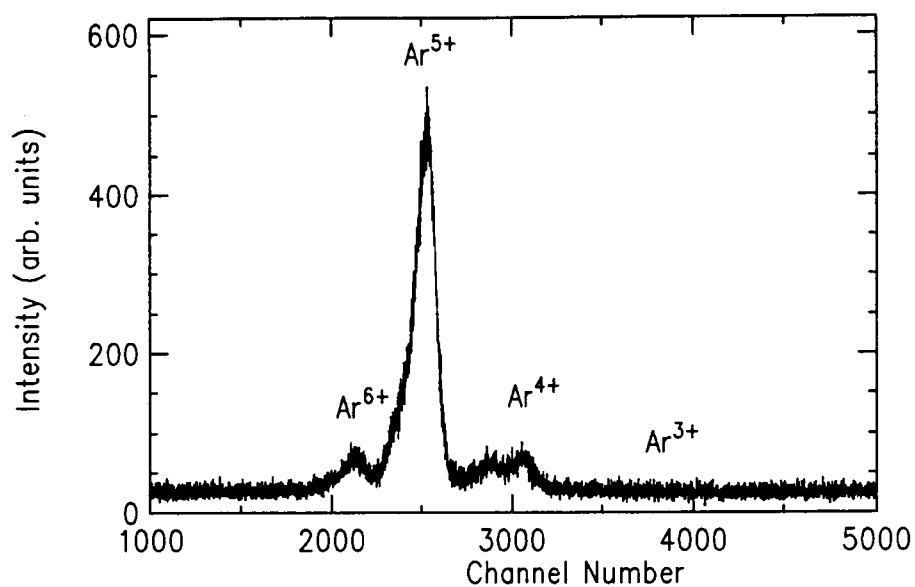


Figure 52: Ion time-of-flight spectrum observed in coincidence with KL_1L_{23} Auger electrons, for 3213.5 eV photon energy.

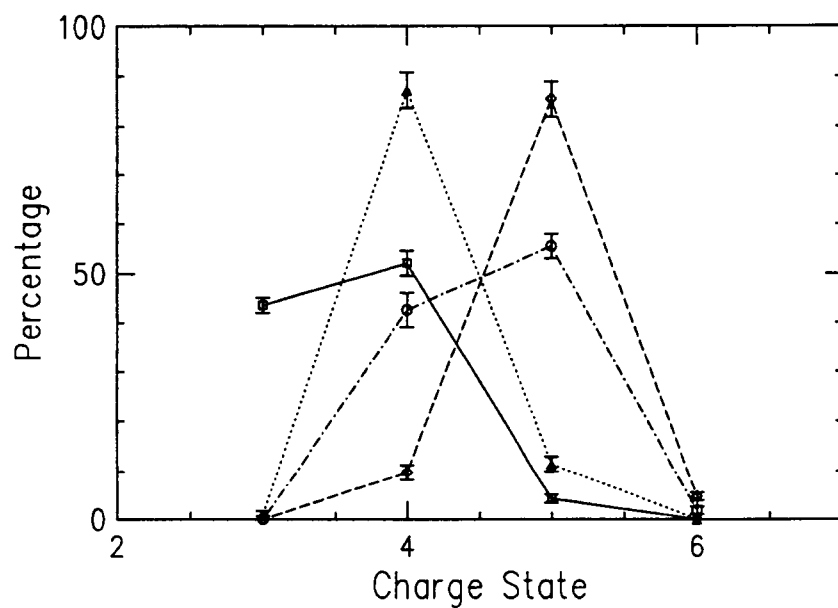


Figure 53: Relative ion-charge-state fractions in coincidence with: (□) $KL_{23}L_{23}$ electrons at resonance, (Δ) $KL_{23}L_{23}$ electrons 10eV above resonance, (○) KL_1L_{23} electrons at resonance and (◇) KL_1L_{23} electrons 10eV above resonance.

From Fig. 53 we can obtain the ion-charge-state fractions at a photon energy above the argon K -edge ($h\nu = 3213.5\text{eV}$); the values are shown in Table 11. From Table 11 we can obtain the branching ratios for the processes such that their first step is a $KL_{23}L_{23}$ or a KL_1L_{23} Auger electron leading to ion charge states $q = 3, 4, 5, 6$. We first analyze the most probable paths starting with a first step $KL_{23}L_{23}$ Auger electron, as shown in Table 12.

The sum of the branching ratios of the two processes that lead to Ar^{3+} is equal to 1.18%. Estimating the fluorescence decay path at only 0.022% (Carlson *et al.* [55]), then $1.158 \pm 0.5\%$ of the Ar^{3+} comes from the decay where two $L_{23}MM$ Auger

Table 11: Ion-charge-state fractions observed in coincidence with $KL_{23}L_{23}$ and KL_1L_{23} Auger electrons at $h\nu = 3213.5\text{eV}$.

	$KL_{23}L_{23}$	KL_1L_{23}
	Rel. Int. %	Rel. Int. %
Ar^{3+}	1.18 ± 0.5	
Ar^{4+}	87.26 ± 3.5	9.84 ± 1.5
Ar^{5+}	11.56 ± 1.5	85.44 ± 3.5
Ar^{6+}		4.72 ± 0.8

Table 12: Decay paths for a first step $KL_{23}L_{23}$ Auger electron.

$KL_{23}L_{23}$	$\longrightarrow$	$L_{23}MM + h\nu[M]$	$\longrightarrow$	Ar^{3+}
		$2(L_{23}MM)n_p'$	$\longrightarrow$	Ar^{3+}
	$\longrightarrow$	$2(L_{23}MM)$	$\longrightarrow$	Ar^{4+}
		$(2(L_{23}MM) + MM)n_p'$	$\longrightarrow$	Ar^{4+}
	$\longrightarrow$	$2(L_{23}M_1M_1) + M_1M_{23}$	$\longrightarrow$	Ar^{5+}

Table 13: Decay paths for a first step KL_1L_{23} Auger electron.

KL_1L_{23}	$\longrightarrow L_{23}MM + L_1MM$	$\longrightarrow Ar^{4+}$
	$L_{23}MM + L_1MM \rightarrow (M_1M_{23})n'_p$	$\longrightarrow Ar^{4+}$
	$\longrightarrow L_{23}MM + L_1L_{23}M \rightarrow L_{23}MM$	$\longrightarrow Ar^{5+}$
	$L_{23}MM + L_1L_{23}M \rightarrow (L_{23}MM + M_1M_{23})n'_p$	$\longrightarrow Ar^{5+}$
	$\longrightarrow L_{23}MM + L_1L_{23}M \rightarrow L_{23}MM + M_1M_{23}$	$\longrightarrow Ar^{6+}$

electrons are emitted, while one electron remains in a n'_p state. From calculations done by Armen *et al.* [21, 56] at 8eV above the Ar K -threshold (3206.3eV), the relative yield for $Ar^{3+}/(Ar^{3+} + Ar^{4+} + Ar^{5+})$ is about 1.78%, i.e. around 1.5 times bigger than the measured value. This decay path is possible if there is a post collision interaction (PCI) between the photoelectron and the KLL Auger electron, in which the slowly moving 1s photoelectron near threshold is recaptured to a n'_p state, with the emission of two $L_{23}MM$ Auger electrons after the initial 1s hole creation. The branching ratio that leads to Ar^{4+} ions decaying by the emission of two $L_{23}MM$ Auger electrons and two $L_{23}MM$ Auger electrons plus a shake-off M_1M_{23} electron recaptured to a n'_p excited state is $87.26\% \pm 3.5\%$. Only one process contributes to the formation of Ar^{5+} where the emission of two $L_{23}MM$ Auger electrons is followed by the shake-off of a M_1M_{23} electron with a branching ratio of $11.57\% \pm 1.5\%$ (only 0.2% of the Ar^{5+} will be in an excited state [56]). From the calculations done by Armen *et al.* [21] the relative yield for $Ar^{5+}/(Ar^{3+} + Ar^{4+} + Ar^{5+})$ is about 8% for photon energies at 8eV above the Ar K -threshold. The most probable decay paths for a first step KL_1L_{23} Auger electron are shown in Table 13.

The branching ratio for the emission of $L_{23}MM + L_1MM$ Auger electrons, underlining the possible shake-off of a M_1M_{23} electron recaptured to a n'_p state is equal to $9.84\% \pm 1.5\%$. These two processes lead to an Ar^{4+} ion final state. The branching ratio for the two processes that lead to Ar^{5+} is $85.44\% \pm 3.5\%$. The main path is a

$L_{23}MM + L_1L_{23}M \rightarrow L_{23}MM$ Auger decay; a weaker path is the emission of the second $L_{23}MM$ Auger electron, followed by a shake-off M_1M_{23} electron recaptured to a n'_p state. A third decay path is the one leading to Ar^{6+} , where the $L_{23}MM + L_1L_{23}M$ Auger electron emission is followed by a $L_{23}M_1M + M_1M_{23}$, Auger electron with a shake-off electron from the M shell. The branching ratio for this process is measured to be $4.72\% \pm 0.8\%$.

5 Conclusions

5.1 Doubly Excited States of Helium Produced by Fast-Ion Impact

From the data presented we infer that, for the carbon energies and charge states employed ($q=4,5,6$), the $2p^2(^1D)$ and the $2s2p(^1P)$ states are primarily populated by first-order mechanisms. For projectiles with higher incident charge state, theory [10] predicts that second-order mechanisms become more important for the production of the $2p^2(^1D)$ state. The behavior of the angular dependence of the singly differential cross section for the production of the $2p^2(^1D)$ state is similar for all three charge states. We can also say that the differential cross section for the production of the $2p^2(^1D)$ state is approximately 1.85 times larger than that for the production of the $2s2p(^1P)$ state in the measured angular range for C^{4+} , and approximately 2.03 times larger for C^{5+} and C^{6+} ions. The ratio of the total double-excitation cross sections for the production of the $2p^2(^1D)$ state to that for the production of the $2s2p(^1P)$ states is approximately equal to 1.97 for $q = 4, 5, 6$. Theoretical calculations by Fritsch *et al.* for C^{6+} at $1.5\text{ MeV}/u$ give a ratio of 0.96, which is about a factor of 2 smaller.

New calculations for the doubly differential cross section (DDCS) are available [49], allowing the direct comparison of the measured spectra with theory. The calculated and measured DDCS are in good agreement concerning the shape and positions of the resonances. There still remains some discrepancy in the absolute value of the DDCS due to the direct ionization component, which are mainly ascribed to the data normalization.

5.2 Cascade-Decay Processes Produced in Photon-Atom Interactions

We have accumulated enough intensity to see a non-isotropic angular correlation between a distinct slice (2) of the LMM Auger group and the precursor $KL_{23}L_{23}$ Auger electrons while the other slices ((1) and (3)) of the subsequent emitted LMM group show an isotropic angular distribution. The non-isotropic angular correlation observed for the distinct LMM slice accompanies the detection of the first-step $KL_{23}L_{23}$ Auger electron in a particular direction, thereby breaking the symmetry so as to reveal the correlation. As the $KL_{23}L_{23}$ and the LMM Auger-electron energies are so different, we believe that the large non-isotropic angular correlation observed represents an alignment effect (of the atomic system) rather than a dynamical post-collision-interaction effect. From a previous assignment [51] of the LMM Auger lines we can see that the main lines contributing to slice (1) are $L_{23}M_1M_1$ transitions; to slice (2) $L_{23}M_1M_{23}$ transitions, and to slice (3) $L_{23}M_{23}M_{23}$ transitions. To better compare the calculations with the experimental results, improved energy resolution needs to be achieved for the LMM Auger group. Because of the small value of the coincidence event rate for this experiment we were not able to run the electron TOF spectrometer at its highest energy resolution. A higher-flux photon beam can compensate the transmission drop in the TOF spectrometer when working at a higher energy resolution, allowing the measurement of the individual LMM Auger lines in coincidence with a given first step KLL Auger electron.

5.3 Photoion Charge-State Distribution and Electron-Ion Coincidences

We studied the importance of decisive experimental parameters to significantly improve the accuracy of the photoion charge state fractions produced by photoionization and photoexcitation of argon atoms. The largest effects on the ion-charge-state

fractions were seen for the lowest charge states, Ar^{1+} and Ar^{2+} , and were ascribed to improved control over the potential difference across the MCP detector and over the CFD threshold.

After determining the optimal values for the most critical experimental parameters, non-coincidence measurement of the photoion charge state fractions were performed as a function of the photon energy. Comparison with previous data [18, 31, 17] shows the important role that the current determination of these parameters has in determining more accurate photoion charge state fractions. We were able to obtain several branching ratios for some decay paths of the cascade-decay process. Coincidence measurements with the first step $KL_{23}L_{23}$ and KL_1L_{23} Auger electrons allowed us to select a given decay path giving us the possibility to obtain the branching ratios of more selected decay processes.

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Appendices

A Details About the Electronics Setup

More detailed diagrams of the electronics used in the coincidence and single spectra are shown in Fig. 54, 55 and 56. Fig. 54 is a diagram of the electron-electron coincidence measurements done at NSLS, where the electron TOF spectrometer and the CMA were used. For these coincidence measurements, the sum signal of the image-particle-detector module or "black box" is used to "gate" the stop signal of the TDC given by the bunch marker signal of the X-ray storage ring. As the sum signal has a jitter of $50nsec$ it can not be used as a timing signal to stop the TDC, because it spoils the time resolution, and therefore the energy resolution of about $1nsec$ of the TOF electronics. Since the sum signal is a rather short pulse, we generated the 'gate' signal with a shape and delay generator having a width of $550nsec$, which is slightly shorter than the distance between the bunch marker pulses of the X-ray storage ring ($567nsec$). To determine the right timing for the coincidence measurements we simulated an event with a pulse generator. The same event was registered by the PSD electronics $2.6\mu sec$ after the TOF electronics. According to this finding, the TOF signal was delayed $2.4\mu sec$, so that both TOF and PSD signals of an event were within our time window of $640nsec$. In order to check the timing of our coincidence setup, we deliberately have set a wrong delay time and got indeed a spectrum different from the coincidence spectrum. The timing signals for the electron-electron coincidence experiments are shown in Fig. 57.

For the electron-ion coincidence measurements the electronics is setup in the same way but now the stop signal comes from the ion TOF spectrometer instead. The signal from the MCP detector of the ion TOF goes through a circuit that decouples the high voltage applied to the MCP from the ion signal before sending it to the amplifier. This decoupling circuit is shown in Figs. 55 and 56. These figures display the electronics diagrams to obtain the time-of-flight spectra with the MP and the PCAME programs, shown for the cases where the electron and ion TOF

spectrometers are used.

A top view of how the PSD is mounted is shown in Fig. 58, where A, B, C and D are connected to the image-particle-detector amplifier (or black box). The resistive anode, rear, middle and front are connected through the voltage divider circuit shown in Fig. 59 to two power supplies, one connected to the anode input (this is a positive $5kV$ power supply) and the other to a front input to bias the MCP plate first encountered by the incident electron (this is a negative $5kV$ power supply). The values of the resistance and the capacitors in Fig. 59 are shown in Table 14. The names assigned by the manufactures to the PSD as rear, middle and front correspond to:

- 1) front = MCP surface first encountered by the incident electron;
- 2) middle = is the contact surface between the first and the second MCP;
- 3) rear = MCP exit surface of the second MCP.

The x and y axes in Fig. 58 are the ones that correspond to the detector orientation in real space (for all three experiments the y axis corresponds to the beam direction). The connection of the A, B, C, and D contacts of the resistive anode needs to be as shown in Fig. 58, to obtain the right correspondence between the x and y axes in real space and the MP program PC display.

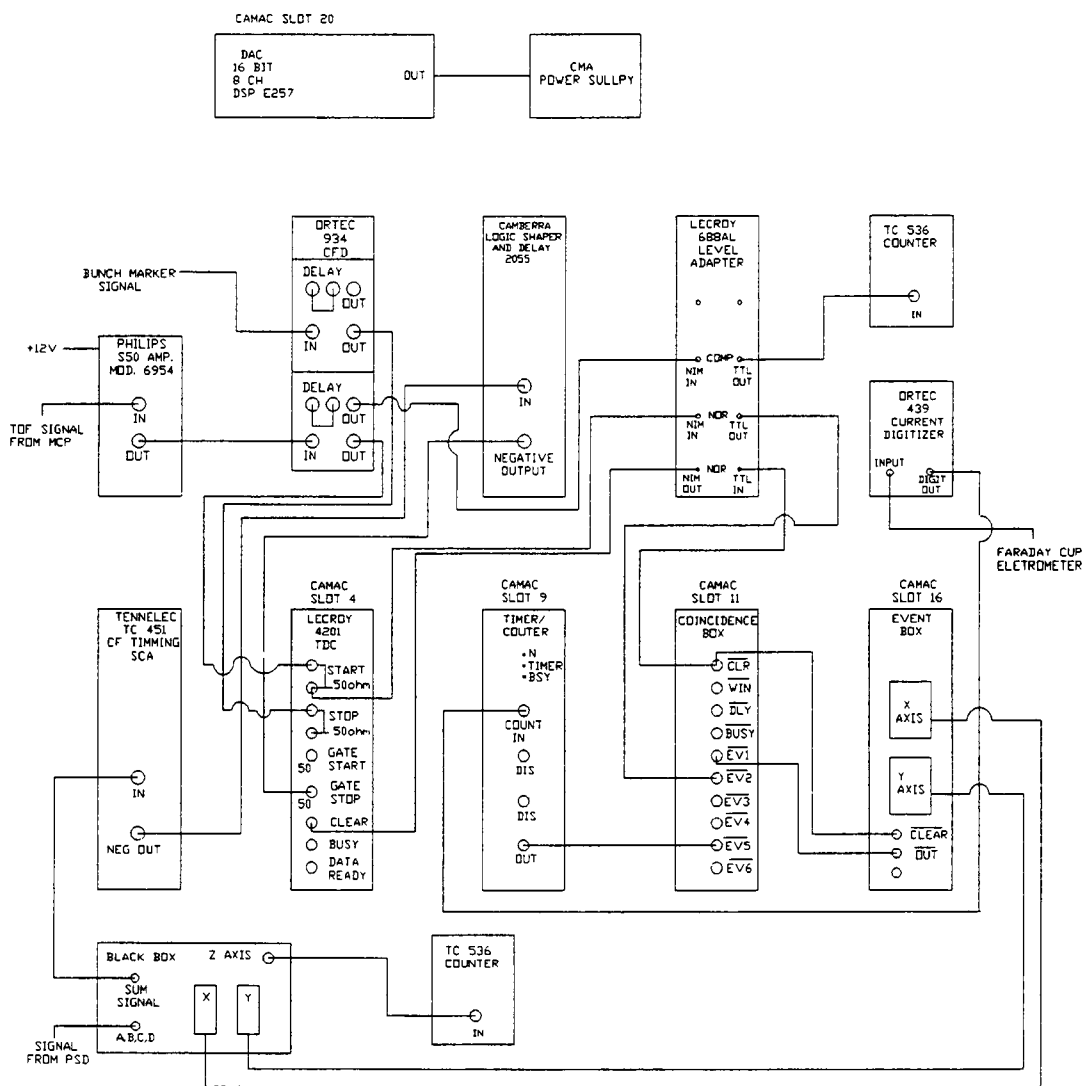


Figure 54: Electronics diagram for electron-electron coincidence measurements.

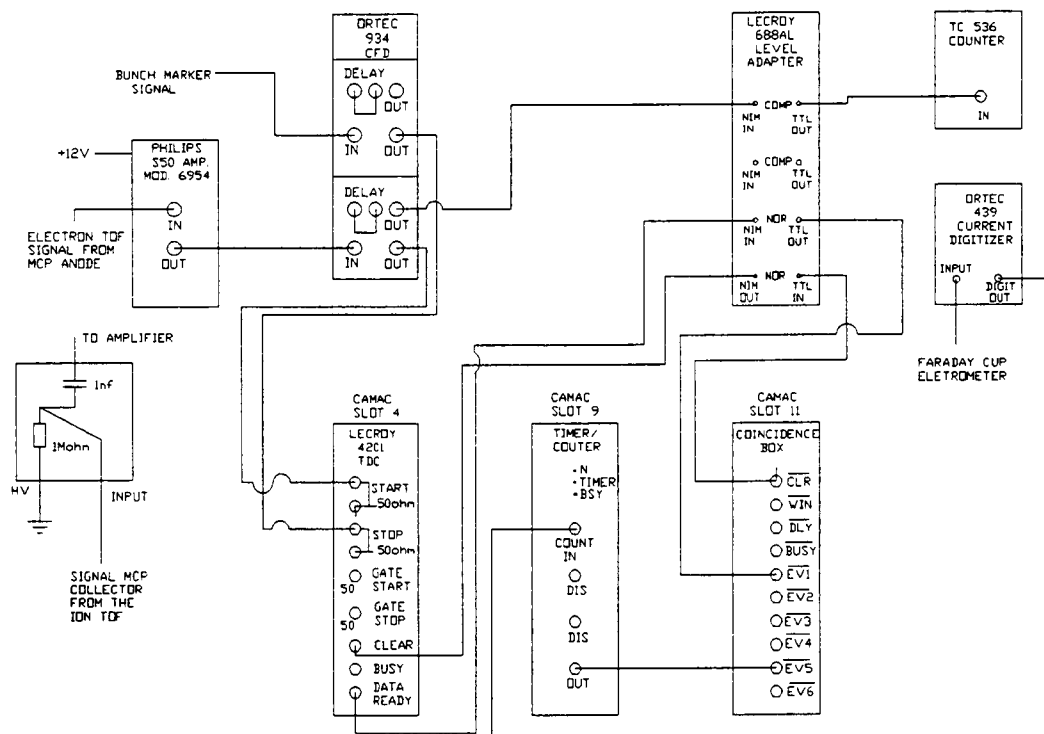


Figure 55: Electronics diagram for singles spectra using the MP program.

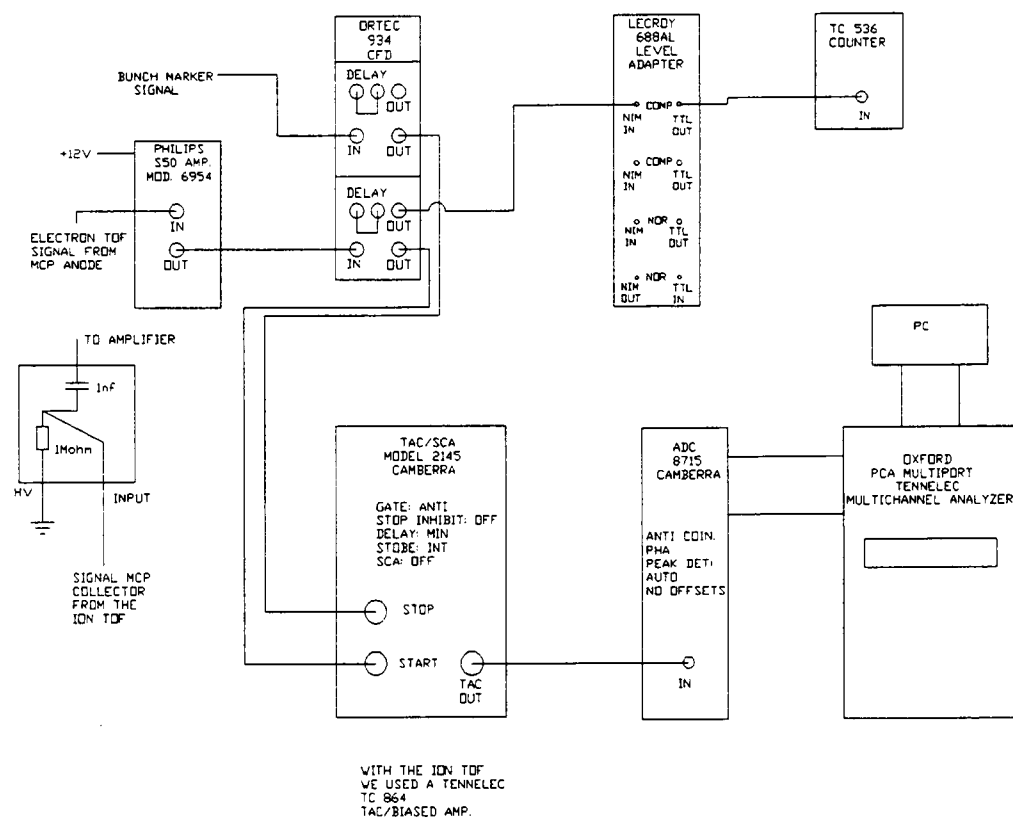


Figure 56: Electronics diagram for singles spectra using the PCAME program.

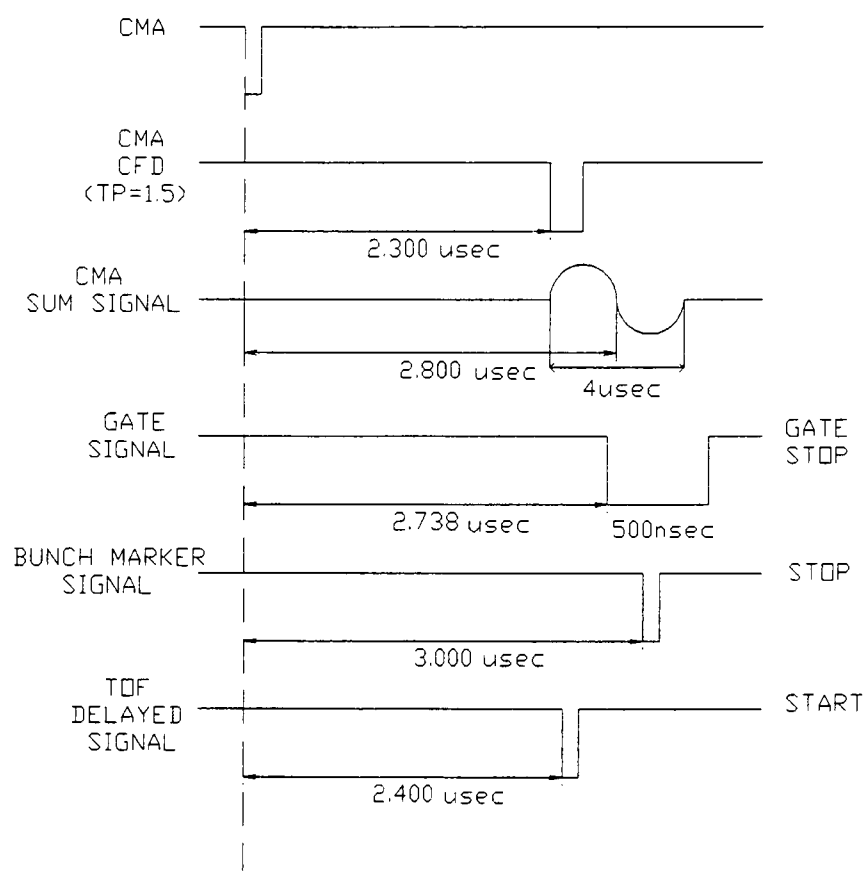


Figure 57: Timing signals diagram for the coincidence setup.

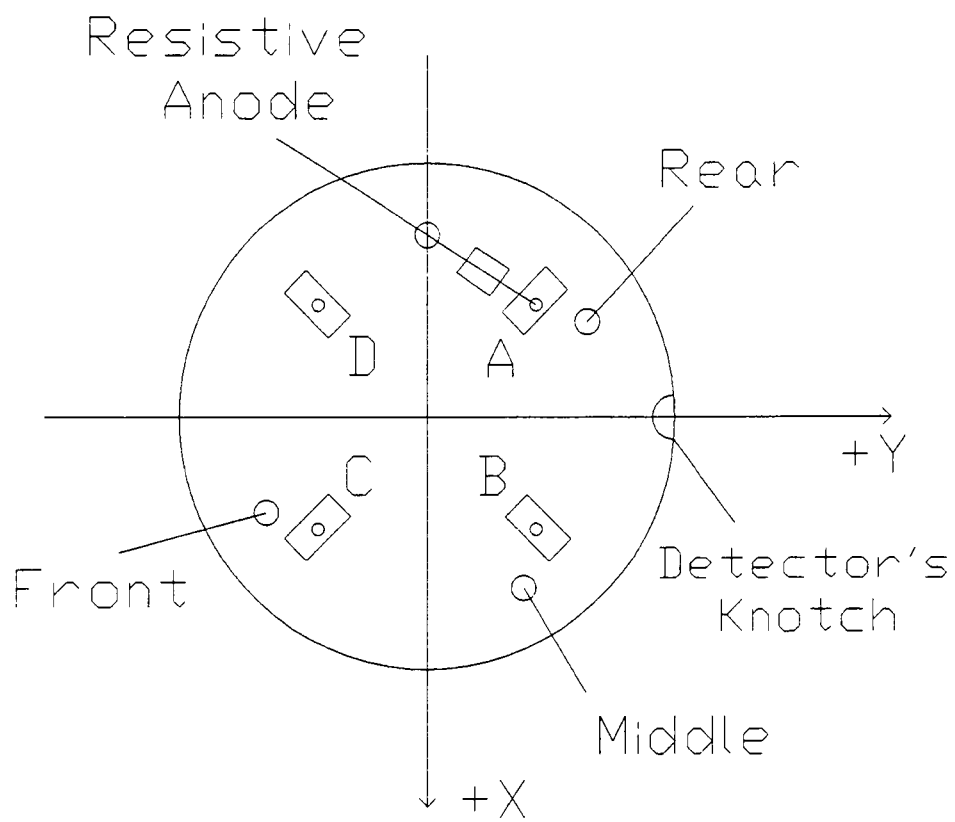


Figure 58: PSD detector connections.

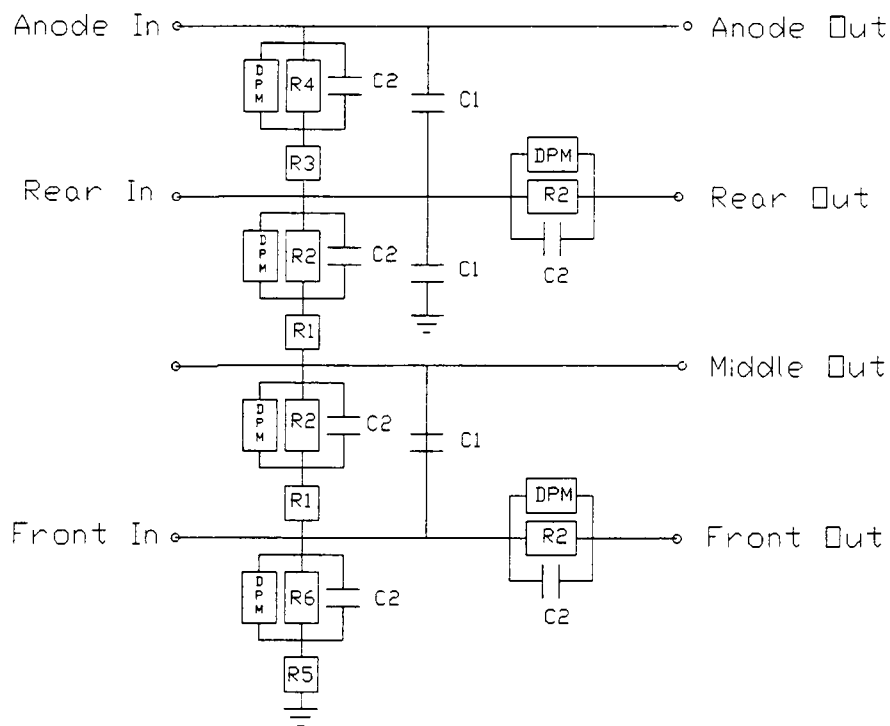


Figure 59: Voltage divider box for the PSD detector.

Table 14: Values for the resistance and capacitors in the voltage divider box.

Resistance	Power	Capacitors
$R_1 = 10M\Omega$	$1/2W$	$C_1 = 0.01\mu F$
$R_2 = 1K\Omega$	$1/10W$	$C_2 = 0.01\mu F$
$R_3 = 3M\Omega$	$1/2W$	
$R_4 = 300\Omega$	$1/10W$	
$R_5 = 20M\Omega$	$1/2W$	
$R_6 = 2K\Omega$	$1/10W$	

B CMA Energy Calibrations for Different Electron Energy Regions

Several energy calibrations of the CMA for different electron-energy regions were conducted. These calibrations are listed below, together with a brief description of how they were done.

1. This first calibration is valid for electron energies between $30 - 40\text{eV}$. It was done using the $N_{45}O_{23}O_{23}$ Xe Auger lines excited by ion impact. The data was taken in UPS mode with a passing energy of 5eV .

$$V_{dac} = \left(\frac{98.6-66.1}{512} \right) ChannelNumber + 66.1$$

$$E(\text{eV}) = 0.02546 ChannelNumber + 26.7859$$

with a standard error for $E(\text{eV}) = 0.1634$

and a standard error for the Channel Number coefficient = 0.0023

2. The second calibration is valid for electron energies between $180 - 210\text{eV}$. It was done with the argon LMM Auger electrons excited by 1KeV electrons. The data was taken in ESCA mode with a passing energy of 10eV .

$$V_{dac} = \frac{51.67-49.23}{255} ChannelNumber + 49.23$$

$$E(\text{eV}) = 0.38507 ChannelNumber + 198.7615$$

with a standard error for $E(\text{eV}) = 0.1675$

and a standard error for the Channel Number coefficient = 0.0014

3. The third calibration is valid for electron energies between $33 - 150\text{eV}$. It was done with the $1s$ argon photoelectrons excited by photons (the argon $1s$ threshold is 3206.3eV). The data was taken in Auger mode (nonretarding mode) with different start and stop DAC voltages for each photon energy.

$$V_{dac} = 0.137775 E(\text{eV}) - 0.505195$$

with a standard error for $V_{dac} = 0.061$

and a standard error for the Energy coefficient = 0.000585.

C Diagram of the Energy Levels of He.

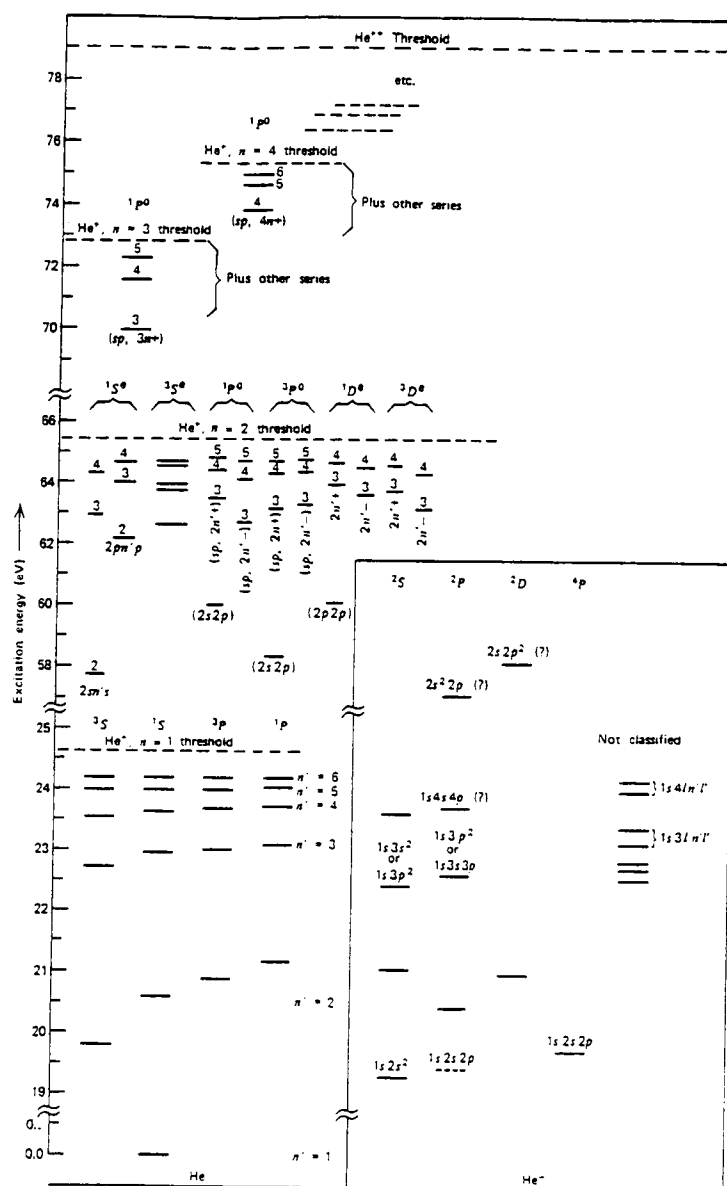


Figure 60: Diagram of the energy levels of He [57].

D Description of MP Program Capabilities.

MP is a program that is used to acquire, display and store singles or coincidence data. From a personal computer (PC) the MP program controls three specially designed and two commercial CAMAC modules via a CAMAC Crate Controller that are part of the data-acquisition system. The program and CAMAC modules were developed by Dr. Leif Liljeby of the Manne Siegbahn Laboratory in Stockholm, Sweden and improved for the kind of experiments that we performed. The program can analyze signals produced by a position sensitive electron detector (PSD) or by nearly any other kind of charged particle detector. It can also be used to control the voltages applied to any electrostatic analyzer plates. The program can either work as a multichannel scaler (MCS) or as a multichannel analyzer (MCA).

In the MCS mode the program makes a step in the channel number after a given time or a given number of counts depending if it is working in timer or counter mode (explained below). In each step the voltages applied to an electrostatic analyzer can be changed to obtain an energy spectrum. Each event is position and energy analyzed by the program.

In the MCA mode the program analyzes the input pulse height and assigns a given channel number for a given pulse height of the input signal. For each mode the program needs a specific definition file (see manual for examples).

For coincidence data the program can work as a MCS and as a MCA at the same time, that is also defined in the definition file. The program monitors slow coincidences between a signal produced by a PSD, and another signal produced by nearly any other kind of detector. The PSD generates a two dimensional position coordinate and an event peak. With the present electronics it is possible to look for coincidences among five different signal, at most; but it was used to perform coincidences between only two signals at a time. As an example of a coincidence measurement; electrons that are selected in energy by an electrostatic analyzer are

later detected in a PSD in coincidence with correlated electrons that are produced jointly in the same interaction region, and are energy analyzed by a time of flight analyzer (TOF) collecting them by a microchannel plate detector (MCP). It is also possible to look for coincidences using any other kind of analyzers as long as only one of them is using a PSD. In the coincidence mode one can measure, for example, electron-electron, electron-ion, ion-ion or electron-photon coincidences (for the moment the system can accommodate one detector that is position sensitive).

The program can also work in list mode, meaning that each event is stored in memory individually so that several spectra and their coincidence relationships can be defined after a given experiment is completed. The data can be sorted back by the program to display and analyze newly defined spectra. This list mode possibility is also defined in the definition file.

A brief description of the three specially designed and two commercial CAMAC modules controlled by the MP program via a CAMAC Crate Controller is given below.

The CAMAC Crate Controller module controls and organizes the CAMAC lines. It accesses all the other modules to read from and write to. It is the connection between the computer and the modules. For more details the manual for the MP program can be consulted.

The Timer/Counter module was specifically built to work with this program. Its main function is to choose between two modes of operation. In the Timer mode the DAC voltage will step after a given time interval set in the program. In the counter mode the DAC voltage will step after the counter has accumulated a given number of digitized Faraday cup counts preselected in the program; normally used for normalization purposes. This module has one output and three inputs. The output is sent to the coincidence module. One of the inputs is called the count input. This input is used in the counter mode. For example when one wants to normalize data with beam current a digitized signal is sent to this input coming

from a current digitizer. There are, in addition, two disable inputs that are not currently used, but can disable the counter with an external signal. All the input signals sent to this module need to be negative true TTL pulses. It also has three LEDs. From top to bottom they are N, time and busy. The N is the slot number in the CAMAC crate, and is equal to nine for this module. The N LED is on when the program, through the controller, access this module. The time LED is on when the timer mode is chosen and off when the counter mode is chosen. The busy LED is on when the counter is enabled.

The coincidence module was also specifically built to use with this program. This module analyzes all the events at its input, processes the information, and tells the computer that an event was analyzed and ready to be stored. From the program one can set a coincidence window, during which time the module waits for all the possible inputs to occur. The module then inform the program that events were present at the inputs during this window. From the program the conversion time can also be set. During this time interval the module processes all the events before the result is sent to the computer through the CAMAC crate controller. This module gives four outputs: window, delay, busy and clear outputs. The window and delay outputs are negative true pulses. Their width is given by the time set in the program. The clear output is a negative true TTL pulse that is present after a given event is processed and the module is ready to accept another event. The busy signal gives a negative true pulse that lasts for the time that the module is busy analyzing a given event. This module has six inputs that need negative true TTL pulses. One is for the timer/counter signal (EV5) and the other five (EV1, EV2, EV3, EV4 and EV6) are available for any event that requires processing, like one coming from the PSD or from the time-of-flight analyzer (TOF). This module has four LEDs. From top to bottom they are N, event enable, LAM enable and abort. The N is a CAMAC function so the LED is on when the CAMAC crate controller accesses this module. For this case N=11. The event enable and LAM enable are commands sent from

the program through the CAMAC Crate Controller to this module. LAM is another CAMAC command and stands for Look at Me (for more details see the CAMAC crate controller manual). When one event is completely analyzed by the computer the program sends the event enable signal so the module is ready for another event. In the same way the LAM enable signal is sent by the computer so the module can set the LAM and the controller can read its bus.

The event module was also specifically built to use with this program. The signal from the position sensitive detector (PSD) is amplified and analyzed by a vendor-provided image-particle-detector module (or "black box" unit). As a result this unit gives the X and Y position of an electron or ion that hit the PDS detector (see the manual for the position sensitive detector for more details). The X and Y position information plus a strobe signal is sent through two wires with a 25-pin connector from this electronic unit to the event module. The event module sends out two signals. One is sent to the coincidence module (EV1) to let it know that there is an event coming from the PSD. A second signal that is only available in the CAMAC bus gives the information to the computer about the position of the event in the PSD. It has a clear input that must be connected to the clear output of the coincidence module. The clear signal must be a negative true TTL pulse. This module has one LED that indicates when an event is there.

The DAC is a digital to analog converter. The digital signal that is sent by the computer for every step of the scan is converted into an analog signal. This signal normally goes to the input of any programmable power supply that controls the voltages in the analyzer (this is a commercial module for what its manual gives more details). The output of this module is a 20-pin connector that connects to a box that has eight outputs (eight BNC connectors). These outputs are 0-10 V analog signals. These values are determined by the Dac 1 Start and the Dac 1 Stop values chosen from the program. These outputs may be connected to different programmable power supplies.

E List of Symbols and Abbreviations.

CMA = Cylindrical mirror analyzer.

He = Helium.

Ar = Argon.

IEM = Independent electron model.

HSCC = Hyperspherical close-coupling.

CC= Close coupling.

ESSR= Electron spectrometry using synchrotron radiation.

SR = Synchrotron radiation.

PCI = Post collision interactions.

MCP = Microchannel plate.

PSD = Position sensitive detector.

DAC = Digital to analog converter.

CAMAC = Computer Automated Measurement and Control; is a standard modular instrumentation and digital interface system.

MP = Name of acquisition program.

PC = Personal computer.

ESCA = Electron spectroscopy for chemical analysis.

UPS = Ultraviolet photoelectron spectroscopy.

HV = High voltage.

TTL = TTL pulse; $+5V$ and $500nsec$ pulse.

NIM = NIM pulse; $-1.5V$ and $\approx 1nsec$ pulse.

TOF = Time of flight.

NSLS = National Synchrotron Light Source.

ORNL = Oak Ridge National Laboratory.

TDC = Time to digital converter.

BM = Bunch marker.

TAC = Time to amplitude converter.

CFD = Constant fraction discriminator.

DDCS = Double differential cross section.

PCAME = The Oxford Instruments Tennelec/Nucleus PCA-MULTIPOINT multi-channel analyzer acquisition system.

MCS = Multichannel Scaler.

MCA = Multichannel Analyzer.

LED = Light-emitting diode.

LAM = Look at me; CAMAC command.

Vita

Leticia S. Pibida was born in Buenos Aires, Argentina. She received her Licenciata (Masters) Degree in Physics from The University of Buenos Aires, Argentina in July, 1991. Her thesis title was Radiative Electron Capture in Atomic Collisions under the supervision of Dr. Ignacio Nemirovsky. In August 1991 she entered The Graduate School of The University of Tennessee, Knoxville. Since August 1992 she has worked on the Doctor of Philosophy degree with a major in Physics under the supervision of Professor Ivan A. Sellin, within the University of Tennessee Atomic Physics Group. The major emphasis of her work has been the study of ion-atom collisions and the interaction of synchrotron radiation with atoms.

The author is a member of the Sigma Pi Sigma honor society and received the 1996 Outstanding Graduate Student Award from The University of Tennessee Department of Physics and Astronomy.