Spin Asymmetries in Deep-Inelastic Electron-Nucleon Scattering — Selected HERMES Results

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Abstract

The HERMES collaboration has collected many millions of deep-inelastic lepton-nucleon scattering events during the years 1995-2000, i.e. the first phase of the experiment. Longitudinally polarised electron or positron beams in the HERA electron storage ring were incident on longitudinally polarised internal atomic gas targets as well as several nuclear gas targets. The primary goal of the HERMES experiment is the study of the spin structure of the nucleon. High precision measurements of double-spin asymmetries in inclusive and semi-inclusive scattering from undiluted polarised atomic hydrogen, deuterium and $^3$He gas targets are presented. These data represent the world’s most precise experimental determination to date of the separate contributions of the spin of up, down and strange quarks to the spin of the nucleon, and the first direct indication of a positive gluon polarisation. The observation of single-spin asymmetries in the azimuthal distributions of hadrons in semi-inclusive deep-inelastic scattering from longitudinally polarised targets is also presented. These single-spin asymmetries can be related to transversity, the only unmeasured of the three leading twist parton distribution functions. Moreover, first measurements of single-spin asymmetries in hard exclusive production of real photons and of pions are reported. Such data can be interpreted in the framework of generalised parton distributions, which is also used when discussing new data on hard exclusive electro-production of the vector mesons $\rho$, $\omega$ and $\phi$. Finally some results from measurement of longitudinal spin transfer in electro-production of $\Lambda$ hyperons, and of transverse $\Lambda$ polarisation in quasi-real photo-production are discussed.
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1 Introduction

HERMES is one of the four experiments at the HERA electron proton collider at DESY. It uses the high current longitudinally polarised electron beam of HERA with a beam energy of about 27.5 GeV together with polarised and unpolarised gas targets internal to the storage ring. Scattered electrons and particles produced in the deep-inelastic electron-nucleon interactions are detected and identified by an open-geometry forward spectrometer with large momentum and solid angle acceptance. The primary scientific goal of HERMES was and still is the detailed investigation of the spin structure of the nucleon. But the physics reach of the experiment extends well beyond this specific aspect of hadronic physics and the experiment can be considered as a facility to explore many details of hadron structure, hadron production and hadronic interactions with electro-magnetic probes at centre-of-mass energies $\sqrt{s}$ of around 7 GeV.

HERMES is based on two novel techniques: longitudinal electron polarisation in a high energy storage ring, which is achieved by a system of spin rotator magnets, and a storage cell target where the polarised atoms from a high intensity polarised source are present as pure atomic species without dilution from unpolarised target material.

The idea to develop an experiment such as HERMES was first introduced at the end of 1987 after the observation of the EMC experiment that only a small fraction of the nucleon’s spin can be attributed to the spins of one of its constituents, the quarks, which became known as the ‘spin-crisis’. Groups from the United States and Germany, with high expertise in spin physics, polarised atomic sources and storage cell technique submitted two letters of intent to DESY, proposing the detailed measurement of the spin-dependent structure functions of the neutron and the proton at HERA. The two groups joined their efforts and the HERMES collaboration was formed, which today consists of 31 groups from 12 countries. At this time HERA was still under construction and it was rather unclear whether high beam polarisation could be achieved at all at such a large storage ring. Only after it had been demonstrated experimentally in 1992 that high transverse electron beam polarisation of about 60% could be routinely obtained at HERA and that internal storage cell targets could be operated reliably in other storage rings, approval was granted in October 1992, pending funding. HERMES was fully approved in July 1993 after also funding of all components was secured. The whole detector and its infrastructure was constructed within about 1.5 years. It was commissioned in early summer 1995.

The first phase of HERMES data taking covers the years 1995 to 2000. Data were taken with longitudinally polarised $^3$He (1995), hydrogen (1996-1997) and deuterium (1998-2000) atomic gas targets, but also with several other unpolarised targets ($^2$H, $^2$D, $^3$He, $^4$He, $^2$N, Ne and Kr). The second phase of HERMES data taking, which will cover the years 2002 to 2006, has just started. The physics program will concentrate on measurements with a transversely polarised hydrogen target for a first determination of the so-called transversity distribution and measurements with high density unpolarised hydrogen targets to study in detail exclusive processes.

In this review first some aspects about the HERA polarised electron beam, the high intensity sources of polarised atoms, the storage cell target, and the HERMES spectrometer are briefly summarised. Then the highlights of the results obtained by HERMES in the period 1995-2000 are presented with special emphasis on results connected to double-spin asymmetries, obtained with both beam and target longitudinally polarised, and single spin asymmetries where only beam or target was polarised, or polarisation was observed in the final state.

The primary goal of HERMES has been to perform measurements that dissect the spin structure of the nucleon, elucidating the role of the various quark flavours as well as that of the gluonic field. These basic constituents are capable of contributing to the nucleon’s spin both via their intrinsic spin as well as via their orbital motion:

$$\frac{1}{2} = \frac{1}{2} \Delta \Sigma + L_z^q + \Delta g + L_z^g,$$

(1)

where $\Delta \Sigma$ denotes the contribution of the quark spins, $\Delta g$ the contribution from gluon spin, and $L_z^q$ and $L_z^g$ are the contributions from the orbital angular momenta of quarks and gluons.

Detailed information about $\Delta \Sigma$ and its decomposition into the contributions from quarks and antiquarks of different flavours can be obtained from double-spin asymmetries of cross sections in inclusive
polarised deep-inelastic scattering, where only the scattered electron is observed, and semi-inclusive deep-inelastic scattering where only one of the leading hadrons is observed. The formalism of polarised deep-inelastic lepton-nucleon scattering is briefly summarised, followed by a detailed discussion of the HERMES results for the polarised structure function $g_1$ for the proton, the neutron and the deuteron, and of the semi-inclusive asymmetries, from which the polarisations of up and down quarks and of up, down, and strange anti-quarks have been determined. In this context also some results on pion multiplicities and the measurement of the flavour asymmetry of the light quark sea are presented.

The polarised gluon distribution is difficult to access. One way to probe it is by a measurement of the dependence of the polarised structure functions on the squared four-momentum transfer $-Q^2$. However, the presently available data are not of sufficient quality and do not cover a large enough kinematic range to seriously constrain the polarised gluon distribution, although there is some indication for its integral to be positive. HERMES has observed a negative asymmetry in the production of pairs of hadrons with high transverse momenta which can be considered as the first experimental evidence for a positive gluon polarisation.

While the measurements of the inclusive and semi-inclusive double-spin asymmetries were planned in detail at the time the proposal was written, there is a substantial fraction of exciting new physics topics and pioneering measurements which were not foreseen at that time, but were luckily contained in the data taken and uncovered during the course of the data analysis. These are on the one hand the first observations of single-spin azimuthal asymmetries in the exclusive production of pions or real photons (deeply virtual Compton scattering), and on the other hand measurements of cross sections and decay angular distributions in exclusive vector meson electro-production, which all can be related to the recently introduced generalised parton distributions, which allow for a unified QCD description of many reactions in electron-nucleon scattering and make it possible to relate them in the most comprehensive way. Especially there is hope that from the second moment of combinations of such generalised parton distributions one can determine the total angular momentum of quarks $J^q = \frac{1}{2} \Delta \Sigma + L_z^2$. With the knowledge of $\Delta \Sigma$ it will thus be possible to determine the orbital angular momentum contribution of quarks to the nucleon's spin. After a short introduction into this new and rapidly developing subject, some of these results are presented and discussed which demonstrate the feasibility of such measurements. These data also show the necessity of a high resolution and high luminosity next generation polarised electron-nucleon facility to study these topics in detail.

Another important unforeseen result is the first observation of a single-spin azimuthal asymmetry in semi-inclusive deep-inelastic pion production recorded with a longitudinally polarised target. This asymmetry can likely be related to a hitherto completely unknown set of quark distributions known as transversity, which are required for a complete description of the partonic structure of the nucleon in leading twist in addition to the well known unpolarised and polarised quark distributions. The asymmetry arises from the small transverse component of the target polarisation seen by the exchanged virtual photon and there is good hope that much larger effects will be observed with a transversely polarised target. The investigation of transversity and the related fragmentation functions will be one of the main goals of the second phase of HERMES data taking if a transversely polarised target is used.

Finally some results about polarisation effects in electro-production of $\Lambda$ hyperons are presented: the investigation of longitudinal spin transfer from the virtual photon to the $\Lambda$ particle and the observation of a surprisingly large positive transverse $\Lambda$ polarisation in quasi-real photo-production.
2 The HERMES experiment at HERA

2.1 The polarised electron beam at HERA

The HERA electron beam has a time structure which allows for up to 220 bunches, which have a length of 27 ps and are separated by 96 ns. This corresponds to a duty factor of $0.3 \cdot 10^{-3}$. The average beam current at injection is up to 50 mA and decreases exponentially with a beam life time of about 10 hours. HERA runs with both electrons and positrons and in fact most of the HERMES data have been taken with a positron beam. For simplicity throughout the whole paper electrons are used for both lepton species.

In high-energy storage rings, electron beams can become transversely polarised through the emission of spin-flip synchrotron radiation [1] in the arcs, the so-called Sokolov-Ternov mechanism. This process involves a small asymmetry in spin-flip amplitudes, which enhances the population of the spin state anti-parallel to the magnetic bending field. This polarisation develops in time according to

$$P(t) = P_m(1 - e^{-t/\tau}),$$

where the asymptotic polarisation $P_m$ and the time constant $\tau$ are characteristics of the storage ring. For an ideal machine without any depolarising effects, the maximum asymptotic polarisation theoretically achievable is $P_m = 0.924$. The rise-time constant $\tau$ increases with the third power of the bending radius $\rho$ of the storage ring and decreases with the fifth power of the beam energy $E$. For the (ideal) HERA storage ring, operated at an energy $E = 27.5$ GeV, the theoretically expected value $\tau_{th}$ is 37 min.

In a realistic storage ring depolarising effects can, however, substantially reduce the maximum achievable polarisation. The depolarising effects also affect the actual rise time, which scales with $P_m$ according to

$$\tau = \left(\frac{P_m}{P_{th}}\right)\tau_{th}.$$  \hspace{1cm} (3)

Thus for a typical asymptotic beam polarisation $P_m = 0.55$, the rise time is about 22 min. Precise alignment of the machine quadrupoles and fine tuning of the orbit parameters is needed to achieve high polarisation. These intricate effects can not be controlled precisely, making it necessary to continuously measure the beam polarisation.

Longitudinal beam polarisation is required for polarised electron scattering experiments at high energies, as for a transversely polarised beam all effects are suppressed by $m_e/E$, where $m_e$ is the electron mass. This is obtained by a pair of spin rotators [2] located upstream and downstream of the HERMES experiment in the HERA East straight section, consisting of interleaved horizontal and
vertical bending magnets. The first spin rotator rotates the spins into the beam direction and the second one turns them back to the vertical direction before the beam enters the next arc (see Fig. 1). During the high-luminosity upgrade of HERA in 2000/2001 additional pairs of spin rotators have been installed around the HERA North and South straight sections to generate longitudinal polarisation also for the two collider experiments H1 and ZEUS.

![Schematic layout of the Longitudinal Polarimeter and its components.](image)

With this system longitudinal electron polarisation has been achieved for the first time in 1994 in a high energy storage ring [3]. Since then values of the equilibrium polarisation in the range 40% to 65% are routinely obtained under normal running conditions. In principle the sign of the beam polarisation can be reversed for every fill which requires moving the magnets of the spin rotators. In practice, a single polarisation direction was used for the measurements in 1995/96. Since then the polarisation direction has been reversed typically every few months. The polarisation of the electron beam is continuously monitored by two polarimeters which utilise the spin dependent cross section for Compton scattering of circularly polarised laser photons on the stored electrons. Transverse beam polarisation leads to a small up-down spatial asymmetry of the back-scattered photons with respect to the orbital plane of the electrons for the two helicity states of the laser beam. Energy and vertical position of individual Compton photons are measured with the Transverse Polarimeter [4] in the HERA West section. For 100% electron polarisation the centre of gravity of the vertical position of the Compton photons on the calorimeter, which is located 65 m behind the interaction point of the laser beam and the electron beam, shifts by about 0.3 mm. Longitudinal beam polarisation modifies the energy dependence of the cross section. The corresponding asymmetry in energy deposition of the backscattered photons for the two helicity states of the laser photon beam is measured with the Longitudinal Polarimeter [5] in the HERA East section, downstream of the HERMES experiment. This polarimeter normally works in the multi-photon mode: it measures the total energy deposited in the detector by $10^3 - 10^4$ Compton photons per bunch and also allows to measure the polarisation of individual bunches. The schematic layout of the Longitudinal Polarimeter is shown in Fig. 2.

A measurement of the beam polarisation to an absolute statistical accuracy of 0.01 requires typically one minute, the fractional systematic uncertainty of the Transverse Polarimeter is < 2%, and that of the Longitudinal Polarimeter 1.6%. A typical rise-time curve, measured simultaneously with both polarimeters is shown in Fig. 3.

### 2.2 The HERMES polarized gas target

The HERMES experiment uses an innovative technique for the polarised target, which is very different from other polarised deep-inelastic scattering experiments, that is a polarised gas target internal to the HERA storage ring. This technique permits essentially background-free measurements from highly polarised nucleons with little or no dilution of the signal from unpolarised nucleons in the target.

A schematic view of the HERMES target region is shown in Fig. 4. Polarised atoms from a polarised
hydrogen (H), deuterium (D) or $^3$He source are injected into an open-ended thin-walled storage cell through which the circulating electron beam of the HERA storage ring passes. The target atoms bounce from wall to wall many times before escaping from one of the cell exits, thus increasing the chance of interaction with the beam. If one furthermore decreases the velocity of the atoms by cooling the cell walls to a temperature of about 65 - 100 K (25 K), the areal target density can be increased by about two orders of magnitude compared to a free atomic beam without cell to values of about $1 - 2 \cdot 10^{14} (1 - 10^{15})$ nucleons/cm² for H and D ($^3$He), where the exact values depend on the dimensions and temperature of the cell chosen and the maximum flux of polarised atoms which can be delivered by the corresponding source. The gas atoms diffuse into the storage ring and are removed by a high speed differential pumping system. Many details of the target system and its components are described in Ref. [6].

The storage cell used for the H and D ($^3$He) measurements was constructed from 99.5 % (99.999 %) pure aluminium, 75 µm (125 µm) thick. The walls are as thin as possible to minimise bremsstrahlung, multiple scattering and energy straggling of particles passing through them. The cell is 400 mm in length, the cross section is elliptical to match the HERA electron beam shape, 29.8 mm wide and 9.8 mm high. For the data taking in 2000 the dimensions were reduced to 21.0 mm by 8.9 mm to increase the target density. Inside the storage cell the gas density exhibits a triangular profile along the beam path.

The cell is coated with Drifilm [7] to avoid depolarisation of the atoms due to wall bounces. In addition during operation a thin ice film which has excellent polarisation preserving properties is produced.
due to a small admixture of oxygen to the hydrogen or deuterium at the dissociator.

A system of collimators in front of the target shields the cell from synchrotron radiation emitted by the electron beam.

The beam wake fields can heat the target cell. This is minimised by avoiding any sharp steps in the beamline profile by a metal mesh, mounted both upstream and downstream of the target providing a smooth transition from the elliptical cell cross section to the circular beam pipe.

A strong and uniform (±1.5%) longitudinal magnetic field of 0.335 T in the target region, provided by a superconducting magnet, decouples the electron and nucleon spins, suppressing depolarisation of the nucleons in the target gas due to transient magnetic fields of the high current bunched electron beam of HERA. Such beam induced depolarisation effects have been studied in detail for the hydrogen target [8]. Working conditions have been found between two resonances where beam-induced depolarisation is undetectable at an uncertainty level below 1 %.

Particles scattered into the spectrometer acceptance exit the target chamber through the 'exit window', a thin (0.3 mm) stainless steel foil (see Fig. 4).

2.2.1 The polarised $^3$He target

The $^3$He target [6] was used in the first year of HERMES operation. Polarised $^3$He is regarded as an effective polarised neutron target, as the proton spins largely cancel, and therefore the $^3$He target is very well suited for the measurement of the polarised neutron structure function $g_n$. The method employed to polarise $^3$He atoms involves direct optical pumping of a meta-stable state [9]. If a weak electric discharge is maintained in a low pressure $^3$He gas, a small fraction of the atoms ($\approx 10^{-6}$) will be in the long-lived $^3S_1$ meta-stable state. Circularly polarised light incident upon the sample along a weak applied magnetic field excites transitions between the $^3S_1$ and $^3P_0$ states. Angular momentum is transferred from the pumping light to the meta-stable atoms which become polarised. Transfer of polarisation to the ground-state atoms is achieved through meta-stability-exchange collisions. In order to create an internal target, $^3$He gas flows through a glass pumping cell, where it is polarised, into the target storage cell. The flow was adjusted to provide a target thickness of $1.0 \cdot 10^{15}$ nucleons/cm$^2$, about a factor of three smaller than the source could in principle deliver, to avoid substantial impact of the target on the lifetime of the HERA beam (see Sect. 2.4).

An infrared laser system consisting of a Nd:LAN crystal in a Nd:YAG cavity was used as the source of the 1.083 μm photons required for the $^3S_1 - ^3P_0$ helium transitions. To reverse the polarisation of the target, the sense of circular polarisation of the laser light was reversed.

The longitudinal target field can be much smaller as compared to the field needed for the H, D target, due to the closed electron shell in helium. A field of 3.4 mT was provided by a rectangular pair of Helmholtz coils of dimension 1560 mm by 1200 mm separated by 640 mm.

The nuclear polarisation of the $^3$He atoms was monitored continuously during the experiment both in the pumping cell and in the target cell by an optical technique [10] involving the observation of light emitted in the de-excitation of electronic states originally excited by the discharge or by the beam pulse passing the target cell [11]. The maximum polarisation observed in the target was 54 %, while the average polarisation over the data taking period was 46 %, with a relative accuracy of 3.4 %.

2.2.2 The polarised hydrogen and deuterium target

A schematic view of the the polarised hydrogen and deuterium target is shown in Fig. 5. It consists of three parts: an atomic beam source (ABS), an atomic Breit-Rabi beam polarimeter (BRP) and a target gas analyser (TGA). The ABS is a conventional atomic beam source [12]. Atomic hydrogen or deuterium is produced in a high frequency or microwave dissociator [13]. An intense beam of small divergence is formed by adiabatic gas expansion through a cooled nozzle and collimators into the vacuum supported by a powerful pumping system. A system of sextupole magnets produces electron spin polarisation states by focusing atoms with electron spin component $+\frac{1}{2}$ and deflecting atoms with electron spin component $-\frac{1}{2}$. The sextupole magnets are high gradient permanent magnets consisting of 24 segments with a maximum pole-tip field of 1.5 T. Nuclear polarisation is obtained by the exchange of hyperfine
occupation using high-frequency transitions. Using various combinations of high-frequency transitions, combinations of specific hyperfine states can be injected into the target, allowing different combinations of electron polarisation and target vector and (in the case of deuterium) tensor polarisation. The atomic beam source provided a highly polarised hydrogen (deuterium) atomic beam with intensities up to $6.5 \cdot 10^{10}$ atoms/s ($5.7 \cdot 10^{10}$ atoms/s) in 2 (3) hyperfine sub-states. Nuclear polarisation values of 0.97 (0.92) at a degree of dissociation of 0.92 (0.95) for H (D) were reached.

The effective polarisation of the target is obtained by measuring the polarisation of a small sample of atoms effusing from a side tube connected to the centre of the cell with an atomic beam polarimeter (BRP) [14]. The axis of both the ABS and the BRP are tilted by 30° with respect to the horizontal plane to ensure that the atoms sampled by the BRP have undergone at least one wall collision and do not stem from the injected beam. The principle of this system is similar to that of the atomic beam source: sextupoles operate as electron spin filters, adiabatic high frequency transitions exchange the population of two or more hyperfine states. A quadrupole mass spectrometer serves as beam detection system. The system is able to measure the population of hyperfine states of the sample effusing out of the storage cell with a systematic error below 0.01 for each hyperfine state. A polarisation measurement with a statistical uncertainty of 0.01 requires 60 seconds.

The target gas analyser (TGA), shown in Fig. 5, measures the relative flux of mass 1 and mass 2 for hydrogen (mass 2 and 4 for deuterium), exiting the sample tube, and thus the fraction of molecules in the target. The TGA is mounted off axis, at an angle of 7°, so as not to interfere with the beam entering the BRP. The atoms and molecules are detected with a quadrupole mass spectrometer.

The polarisation and degree of dissociation of the sample probed by the BRP and the TGA are not identical to the effective values as seen by the electron beam. This arises partially because of recombination of atoms in the target cell and the possibility of finite polarisation of molecules and also because sampling corrections have to be applied. These take into account differences in the number of wall collisions of the atoms within the sample beam and the target cell and of the possibility of a non-uniform surface in the cell and sample tube, and have been determined by extensive Monte Carlo Simulations [14, 15, 16, 17]. The mean absolute polarisation seen by the electron beam ($|P^T|$) for the hydrogen (deuterium) target was $0.853 \pm 0.03 (0.845 \pm 0.027)$.

2.3 The HERMES spectrometer

A schematic side view of the HERMES spectrometer is shown in Fig. 6. The spectrometer is described in detail elsewhere [18]. It is a forward spectrometer with a dipole magnet surrounding the electron and proton beam pipes, providing an integrated field of 1.3 Tm. The magnet gap is divided into two identical sections, "up" and "down", by a horizontal septum plate that shields the electron and proton beams from the spectrometer's magnetic field. Consequently, the spectrometer is constructed as two identical halves,
mounted above and below the beam pipes. Scattered electrons and hadrons produced in the inelastic reactions can be detected and identified within an angular acceptance \( \pm 170 \) mrad horizontally and 40 - 140 mrad vertically. For tracking in each spectrometer half several tracking chambers (microstrip-gas counters, multiwire proportional chambers and drift chambers) before, inside and behind the magnet are used. Due to the 40 cm long extended gas target tracking detectors in front of the magnet are needed for the determination of the interaction vertex, the polar and azimuthal scattering angles of the particles and the initial trajectory for the determination of the particle's momentum. Together with the wall of the storage cell and the exit window of the target chamber these detectors cause multiple scattering and energy straggling and are thus limiting the track resolution of the entire spectrometer. Fast track reconstruction is achieved by a pattern-matching algorithm and momentum look-up method [19]. For electrons with momenta between 3.5 and 27 GeV, the average angular resolution is 0.6 - 0.3 mrad and the average momentum resolution \( \Delta p/p \) was 0.7 - 1.3% aside from bremsstrahlung tails. The resolution deteriorated somewhat, after in 1998 the gas threshold Čerenkov counter was replaced by a Ring Image Čerenkov counter (RICH) with more material in the particle's path. The trigger for scattered electrons is formed by a coincidence of signals from three hodoscope planes (H0, H1 and H2) with those from a radiation-hard lead-glass calorimeter, requiring an energy of greater than 1.4 GeV (for some of the data taking periods 3.5 GeV) to be deposited locally in the calorimeter. Electron-hadron discrimination and particle identification is accomplished using the information from four particle identification detectors: the lead-glass calorimeter, a pre-shower detector consisting of a scintillator hodoscope preceded by two radiation lengths of lead located directly before the calorimeter, a six module transition radiation detector (TRD), and a Čerenkov counter. These detectors serve to suppress the large background of hadrons, mostly pions from photo-production, which in the worst part of the kinematic plane is about 400 times higher than the electron yield. A likelihood method, based on the empirical responses of each of the four detectors is used to discriminate between electrons and hadrons. The system provides electron identification with an average efficiency of 98 - 99% and a hadron contamination of less than 1%. The luminosity of the experiment is measured by detecting pairs of electrons from Møller scattering \( e^-e^- \rightarrow e^-e^- \) of beam electrons off the target gas electrons, or in the case of a positron beam \( e^+e^- \) pairs from Bhabha scattering \( e^+e^- \rightarrow e^+e^- \) or photon pairs from annihilation \( e^+e^- \rightarrow \gamma\gamma \). The scattered particles exit the beam pipe at 7.2 m downstream the target center and are detected in coincidence by a pair of two small very radiation-hard electro-magnetic calorimeters, with a horizontal acceptance of 4.6 to 8.9 mrad, which are mounted symmetrically on either side of the beam line.
The whole spectrometer, apart from the target and the muon wall, is mounted on a movable platform to allow access to the HERA tunnel in case of major repairs or replacement of accelerator components. The total length is restricted to about 8 m and all the detectors are very tightly packed.

Below the main features of the detectors as shown in Fig. 6 are summarised. The quoted numbers of modules refer to one of the detector halves above or below the beam, the number of channels refers to the total number of the whole detector system. The coordinate system used by HERMES has the z axis along the beam direction, the y axis pointing upwards, and the x axis horizontal, pointing towards the outside of the storage rings. All wire chambers have planes with wires oriented either in the vertical (for x measurement), or tilted 30° right or left (U and V planes). X', U' and V' indicate drift chambers planes staggered by half the width of a drift cell with respect to the corresponding X, U and V planes.

- **VC1/2**: These Vertex Chambers [20], which are installed directly behind the exit window of the target chamber, are microstrip gas chambers (MSGC) with a cell width of 0.193 mm. There are two modules with VXY configuration, where the stereo angle for the V plane is 5°, resulting in 24800 channels. The gas composition is DME/Ne (50/50), the resolution/plane (σ) is 65 μm. These chambers were not operational in 1999-2000.

- **DVC**: The Drift Vertex Chambers are immediately following the VCs. They have been installed in 1996 and are in use since the data taking in 1997. They help to increase the redundancy of the tracking system in the front region and have a larger geometrical acceptance than the other detectors extending vertically from ±35 mrad to ±270 mrad and covering ±200 mrad horizontally, mainly to increase the efficiency for the detection of events from charmed particle decays. The width of the drift cell is 6 mm, the module has 6 planes with the configuration XX'UU'VV', the total number of wires is 1088. The non-flammable gas mixture is the same as for all the other drift chambers: Ar/CO2/CF4 (90/5/5), the spatial resolution per plane (σ) is 220 μm.

- **FC1/2**: The Front Chambers [21] provide good spatial resolution immediately in front of the magnet. They are drift chambers with a cell width of 7 mm, each module has six planes with a UU'XX'VV' configuration, resulting in 2304 channels. The measured resolution per plane is 225 μm.

- **MC1/3**: The magnet chambers MC1 through MC3 [22] are multiwire proportional chambers with a spacing of the readout wires of 2 mm and a corresponding resolution per plane of 700 μm. They are located in the gap of the magnet and were originally intended to help resolve multiple tracks in case of high multiplicity events and to improve track reconstruction in case of missing planes in the front region. Since low background and good performance of the front detectors have made this unnecessary, their primary function is now the momentum analysis of relatively low energy particles from the decay of Λ hyperons or K mesons, for example. Each module has three planes of the configuration UXV, the total number of channels is 11008, and the gas mixture used is Ar/CO2/CF4 (65/30/5).

- **BC1-4**: The Back Chambers BC1 through BC4 [23, 24] are drift chambers with a cell width of 15 mm. Each of the four modules has six planes with UU'XX'VV' configuration. The active areas of the chambers are chosen according to their z-position and the acceptance of the spectrometer, resulting in 7680 channels. The measured resolution per plane is in average 275 μm (300 μm) for BC1/2 (BC3/4) with a minimum of about 210 μm (250 μm) in the middle of the drift cell.

- **Čerenkov**: During the data taking periods 1995-1997 the Čerenkov counter was a gas threshold counter filled with N2 or a mixture of N2 and C4F10 at atmospheric pressure. This counter served to distinguish pions from heavier hadrons in the momentum range up to 13.5 GeV, with a pion threshold of 5.5 GeV in 1996 (radiator N2) and below 4 GeV in 1996/97 (radiator N2/C4F10 (70/30)). In 1998 this counter was replaced by a dual-radiator Ring Image Čerenkov counter (RICH) [25], with clear silica aerogel [26] (refractive index n = 1.03, γthres ≃ 4.2) at the front.
of the detector and C$_4$F$_{10}$ ($n = 1.0005$, $\gamma_{\text{thresh}} \simeq 32$) as radiator in the rest of the detector volume. This RICH, which was designed and constructed within 1.5 years only, allows full $\pi$, $K$, $p$ separation essentially for all particle momenta accepted by the spectrometer ($> 1$ GeV).

- **TRD:** The Transition Radiation Detector consists of 6 modules containing each a radiator with plastic fibers of about 20 $\mu$m diameter as radiator material and a Xe/CH$_4$ (90:10) filled proportional chamber with vertical wires separated by 1.27 cm. Both electrons and hadrons deposit energy in the detector due to the ionisation of the chamber gas, but only electrons produce transition radiation in the HERMES energy regime. Combining the information from several modules and using the truncated mean method one obtains an average pion rejection factor of about 150 for 90% electron efficiency.

- **Calorimeter:** The function of the calorimeter [27] is to provide a first level trigger for scattered electrons, based on energy deposition in two adjacent columns, to suppress pions by a factor $\geq 10$ at the first level trigger and $\geq 100$ in the offline analysis, to measure the energy of electrons and photons from radiative processes, Deeply Virtual Compton Scattering or $\pi^0$ and $\eta$ decays. Radiation resistant Fi101 lead-glass [28] blocks with a front area of $9 \times 9$ cm$^2$, and a length of 50 cm (about 18 radiation lengths) are stacked in two 42 x 10 arrays each above and below the beam. The energy response can be parameterised as $\sigma(E)/E = (5.1 \pm 1.1) \times 10^{-2}/\sqrt{E[\text{GeV}]} + (1.5 \pm 0.5) \times 10^{-2}$, the spatial resolution of the impact point is about 0.7 cm and the pion rejection factor about 2500, integrated over all energies in combination with the pre-shower detector, for a 95% electron detection efficiency.

- **Hodoscopes:** The scintillator hodoscopes H1 and H2 provide a fast signal that is combined with the calorimeter to form the first level trigger. Both counters are composed of 42 vertical scintillators with a width of 9.3 cm to match the size of the calorimeter. In addition H2, which has a passive radiator of two radiation lengths of lead in front of it and acts as a pre-shower-counter, provides discrimination between electrons and hadrons with a pion rejection factor of $\sim 10$ and 90% efficiency for electron detection. The hodoscope H0 in front of the magnet consists of a single sheet of scintillator. It was introduced in 1996 to eliminate triggers, caused by showers generated by the proton beam, by distinguishing forward from backward going particles using the time of flight between the front and rear scintillators.

- **Luminosity monitor:** The Luminosity monitor [29] consists of Čerenkov crystals of NaBi(WO$_4$)$_2$, which have a very high radiation hardness of the order of $7 \cdot 10^6$ Gy, a radiation length of 1.03 cm and a Molière radius of 2.38 cm. Each calorimeter consists of 12 crystals of size $22 \times 22 \times 200$ mm$^3$ in a $3 \times 4$ ($x \times y$) array. The energy resolution determined with a $3 \times 3$ matrix is $\sigma(E)/E \approx (9.3 \pm 0.1) \times 10^{-2}/\sqrt{E[\text{GeV}]}$.

- **Muon Wall:** The Muon Wall is a one meter deep iron wall. It is installed behind the experiment at the beginning of the HERA tunnel. Together with the Muon Hodoscope behind it, it serves to enhance the muon identification capability inside the acceptance of the spectrometer magnet and thereby improving the identification of semileptonic decays of charmed $D$ mesons, like $D^0 \rightarrow K^- \mu^+\nu_\mu$. It also serves to shield the spectrometer from background arising from the proton beam.

- **Wide Angle Muon Hodoscopes:** These counters are installed outside the standard HERMES acceptance, one in front of the magnet between the field clamps and the body of the magnet and one directly behind hodoscope H1. They match the increased vertical angular acceptance of the DVC chambers. Their main purpose is to increase the acceptance for charmed particle decays, especially for the detection of muon pairs from the decay $J/\Psi \rightarrow \mu^+\mu^-$, since at HERMES energies the opening angle between these decay products is generally larger than the angular acceptance of the magnet.

- **Silicon:** The main purpose of the so-called 'LambdaWheels' (LW) is to greatly increase the acceptance for the decay products of $\Lambda$ particles. The yield of reconstructed $\Lambda$ hyperons is enlarged
by about a factor of four as compared to $\Lambda$ decays that are identified from fully reconstructed tracks that pass the entire spectrometer. The detector consists of two wheel-shaped layers of double-sided silicon-strip counters, positioned at $z = 45 \text{ cm}$ and $z = 50 \text{ cm}$ within the 'pumping cross' of the target chamber. The outer radius is 17.5 cm, the inner radius 5.4 cm. Hence, there is a small geometrical overlap with the acceptance of the other detector systems of the HERMES spectrometer. Each LW detector consists of twelve 30 degree segments. The strips are oriented parallel to the side faces of each segment, i.e. with a stereo angle of 30°, which makes all strips accessible from the edge at the outside radius, where there is room to place the electronics. The width of the strips is 160 µm, which results in 437 strips per side of each segment and a total number of 20976 readout channels. The 'LambdaWheels' have been fully installed during the long HERA shut-down in 2000/2001, following the successful commissioning of a single LW module in 2000.

2.4 Collected data

The first phase of the HERMES experiment spans the period between early summer 1995 and August 2001. In 1995 the polarised $^3\text{He}$ target was used for the commissioning of the HERMES apparatus and the first data taking. The areal density of the target was restricted to $10^{15}$ nucleons/cm² (about 3 times lower than the performance of the polarised source would have allowed), to keep the impact of the HERMES target on the lifetime of the electron beam small (see below).

In 1996 and 1997 data were taken with the longitudinally polarised atomic hydrogen target with an areal target density of $7 \times 10^{13}$ nucleons/cm², limited by the atomic flux delivered by the atomic beam source (ABS).

Finally in the years 1998 to 2000 data were taken with the longitudinally polarised deuterium target. Most of the data were collected in the year 2000, when HERA performed extremely well, and the luminosity was increased compared to previous years by decreasing the vertical dimensions of the storage cell target and cooling the cell to 65 K. Thereby the areal target density was increased to a value of about $2 \times 10^{14}$ nucleons/cm².

In addition, unpolarised data for several nuclear species ($^2\text{H}$, $^3\text{D}$, $^3\text{He}$, $^4\text{He}$, $^2\text{N}$, $^2\text{Ne}$ and $^9\text{Kr}$) have been collected with higher target density during short dedicated periods in the years 1995-1999. In 2000 very high unpolarised target densities were available for about one hour at the end of each HERA fill.

The accumulated number of polarised deep-inelastic (DIS) events collected in the years 1996-2000 as determined from an online analysis of the data is shown in Fig. 7 as a function of the number of days since the beginning of the run in that year. It can be seen that the year 2000 was particularly successful due to the increased target density, and the fact that HERA ran extremely well. A typical example of the excellent performance is shown in Fig. 8 where the current in the proton ring (upper curves), the exponentially decreasing current in the electron ring (middle curves) and the lifetime of the electron beam (lower curves) are shown for three subsequent fills during 48 hours in July 2000. Note the sharp decrease of electron beam lifetime from a value of more than ten hours to a value between one and three hours and also the much stronger decrease of electron current with time at the end of each fill. During these periods high density unpolarised gas has been injected into the target cell, hereby increasing the beam losses due to bremsstrahlung in the target and thus decreasing the lifetime of the beam. This is related to the target species and the target density by:

$$\frac{1}{\tau(Z)} = \frac{1}{\tau_0} \left( \frac{Z}{A} (Z+1) \ln \frac{183}{Z^{1/3} \cdot 8.1 \times 10^{25} \text{ nucleons cm}^{-2}} \right),$$

where $Z$, $A$ are charge and mass number of the target material, $n$ is the target areal density in nucleons cm⁻², and $\tau_0 = 21 \mu$s is the revolution time around the ring.

For standard data taking in parallel with the collider experiments H1 and ZEUS, the lower lifetime limit due to the HERMES target alone was set to about 45 to 50 hours, to not affect the total lifetime of the beam by more than about 20%. During the so-called end-of-fill runs the beam lifetime was reduced to 1-2 hours. About 25 millions of DIS events were collected with unpolarised targets.
Figure 7: Cumulative number of polarised DIS events collected by the HERMES experiment in the years 1995 - 2000.

Figure 8: Example of the HERA intensity profile of the proton and the electron beam (upper curves) and the lifetime of the electron beam (lower curve) during 48 hours in July 2000. Note the drop in lifetime at the end of each fill, which is caused by the high unpolarised target gas densities used in these periods.
3 Spin-dependent quark distributions and structure functions in QCD

The topic of the spin structure of the nucleon has been summarised in several previous reviews. The reader is especially referred to the very extensive review of Lampe and Reya [30], where all the theoretical aspects have been discussed in detail and the recent review by Filippone and Ji [31], where the actual theoretical and experimental status is presented. In the following section therefore only those variables are introduced that are required in the context of the present review.

3.1 Kinematics, the cross section and asymmetries for spin-dependent DIS

HERMES performs deep-inelastic electron-nucleon scattering (DIS) experiments

\[ e + N \rightarrow e' + X, \]  

where a point-like electron \( e \) of energy \( E \approx 27.5 \text{ GeV} \) is scattered off a nucleon \( N \) (proton, neutron), thereby exciting it to a hadronic final state \( X \) with an invariant mass \( W \) much larger than the nucleon mass \( M \).

In lowest order perturbation theory this interaction proceeds via the exchange of a neutral virtual boson (7, \( Z^0 \)) and can be interpreted in QCD as the incoherent sum of elastic scattering off quarks of any flavour \( q \). At HERMES energies contributions from \( Z^0 \)-exchange can be safely neglected, therefore only the electromagnetic interaction in the approximation of one-photon exchange is discussed in this review. Moreover only the three light quark flavours up (u), down (d) and strange (s) are of relevance at this energy.

\( k = (E, k) \) and \( k' = (E', k') \) are denoted as the four-momenta of the incident and scattered electron, \( q = (\nu, q^3, p = (E, p^J, p_h = (E_h, p_h) \) as those of the exchanged virtual photon, the target nucleon, and a hadron produced in the interaction. This process is depicted in Fig. 9. It can be characterised by the following three Lorentz-invariant quantities

\[ q^2 = -Q^2 = (k - k')^2 \]  

squared four-momentum transfer, \( s = (p + k)^2 \) squared center-of-mass energy, \( W^2 = (p + q)^2 \) squared invariant mass of the photon-nucleon system.

As only fixed target experiments are discussed, \( p = (M, 0, 0) \). Under the given conditions the electron energies \( E, E' \) are much larger than the electron mass \( m_e \) such that this quantity can be safely neglected.

Thus we obtain in the nucleon rest system

\[ Q^2 = 4EE'(1 - \cos \theta), \]  

\[ s = M^2 + 2ME, \]  

\[ W^2 = M^2 + 2M\nu - Q^2. \]
where $\theta$ is the electron scattering angle in the laboratory system, $\nu = E - E'$ the energy of the virtual photon transferred from the electron to the nucleon, and $M$ the nucleon mass. We can also express the process in terms of two dimensionless scaling variables

$$y = \frac{p \cdot q}{p \cdot k} = \frac{\nu}{E},$$

(12)

$$x = \frac{Q^2}{2p \cdot q} = \frac{Q^2}{2M\nu}.$$  

(13)

Here $y (0 \leq y \leq 1)$ is the fractional energy-transfer from the electron to the nucleon, and the Bjorken scaling variable $x (0 \leq x \leq 1)$ is a measure for the inelasticity of the process. For elastic scattering we have $W^2 = M^2$, and consequently $x = 1$, while for inelastic processes $W^2 > M^2$ and $x < 1$. The variable $x$ can also be interpreted as the fraction of the nucleon’s light-cone momentum carried by the struck quark.

In semi-inclusive deep-inelastic scattering (as discussed in section 4.4.1) hadrons produced in the interaction are also observed in addition to the scattered electron. These hadrons are characterised by two quantities, the fractional energy $z$ and the Feynman variable $x_F$

$$z = \frac{p \cdot p_h}{p \cdot q} = \frac{E_h}{\nu}.$$  

(14)

$$x_F = \frac{p^*_h}{|q|} \approx \frac{2p^*_h}{W}.$$  

(15)

In these definitions $E_h$ is the energy of the produced hadron in the nucleon rest system and $p^*_h$ is its longitudinal momentum with respect to the virtual photon direction in the photon-nucleon centre-of-mass system. Hadrons originating from the struck quark can be preferentially found in the so-called current fragmentation region, $x_F > 0$, with high fractional energy $z$. Hadrons originating from the target remnants can be preferentially found in the target fragmentation region, $x_F < 0$, with small fractional energy $z$.

Information about the internal spin structure of the nucleon can be extracted from deep-inelastic scattering experiments in which a polarised nucleon target and a longitudinally polarised charged lepton beam are used (for transverse beam polarisation all effects are suppressed by $m_e/E$). In this case one also has to take into account the orientation of the polarisation vector $\vec{P}_{\Pi}$ of the target nucleon relative to the direction of the incoming lepton $\vec{k}$, denoted by $\alpha$, and the angle $\phi_S$ between the polarisation plane (formed by $\vec{P}$ and $\vec{P}_{\Pi}$) and the scattering plane (formed by $\vec{k}$ and $\vec{k}'$). These quantities are illustrated in Fig. 10.

![Figure 10: Definition of the angles $\alpha$ and $\phi_S$, describing the orientation of the nucleon spin direction $\vec{P}_{\Pi}$.](image)

The spin-dependent part of the deep inelastic cross section, which is obtained by reversing the target
polarisation, is given by [32]:

\[
\frac{d^3\sigma(\alpha) - d^3\sigma(\pi + \alpha)}{dxdydz} = \frac{e^4}{2\pi^2 Q^2} \left\{ \cos \alpha \left[(1 - y/2 - y^2\gamma^2/4)g_1(x, Q^2) - \frac{y}{2}\gamma^2 g_2(x, Q^2)\right] - \sin \alpha \cos \gamma \sqrt{1 - y} \left[\frac{y}{2}g_1(x, Q^2) + g_2(x, Q^2)\right] \right\}
\]

(16)

The kinematical factor \(\gamma\) is defined as \(\gamma = \sqrt{Q^2/\nu} = 2Mx/\sqrt{Q^2}\), and \(g_1(x, Q^2)\) and \(g_2(x, Q^2)\) are the spin-dependent structure functions. These structure functions depend on \(Q^2\) due to elementary QCD processes such as emission (absorption) of gluons by quarks, production (annihilation) of pairs of quarks and anti-quarks from gluons, the coupling of three or four gluons to each other and higher order processes. In the following the \(Q^2\) dependence of the structure functions will be omitted unless it is explicitly needed in the corresponding context.

Experimentally, the spin-dependent structure functions \(g_1(x)\) and \(g_2(x)\) can be determined by combining asymmetry measurements off a longitudinally (\(\alpha = 0^\circ, 180^\circ\)) and a transversely (\(\alpha = 90^\circ, 270^\circ\)) polarised target. For a longitudinally polarised target the cross section is dominated by the contribution of \(g_1(x)\), the term proportional to \(g_2(x)\) being only a small correction. In the transverse case all effects are suppressed by \(\gamma = \sqrt{Q^2/\nu}\) but both structure functions enter the cross section difference with about the same weights. Because \(g_1(x)\) and \(g_2(x)\) enter both in the spin-dependent cross section expression, a precise determination of \(g_1(x)\) from a longitudinally polarised target alone is not possible. At least some knowledge about the magnitude and \(x\) dependence of \(g_2(x)\) is required.

The cross section asymmetries \(A_{||}\) and \(A_{\perp}\) are defined as

\[
A_{||} = \frac{\sigma_{||}^{11} - \sigma_{||}^{11}}{\sigma_{||}^{11} + \sigma_{||}^{11}} \quad \text{and} \quad A_{\perp} = \frac{\sigma_{\perp}^{1\perp} - \sigma_{\perp}^{1\perp}}{\sigma_{\perp}^{1\perp} + \sigma_{\perp}^{1\perp}} ,
\]

(17)

where the arrows indicate the polarisation directions of electron beam and target. Data on \(A_{||}\) and \(A_{\perp}\) can be obtained by reversing the polarisation directions of a longitudinally or a transversely polarised target. The measured lepton asymmetries are related to the longitudinal and transverse asymmetries \(A_1\) and \(A_2\) of the exchanged virtual photon by the relations

\[
A_{||} = D \cdot (A_1 + \eta \cdot A_2) \quad \text{and} \quad A_{\perp} = d \cdot (A_2 - \zeta \cdot A_1) ,
\]

(18)

(19)

The kinematic factors \(D, d\) and \(\eta, \zeta\) entering these expressions are defined by

\[
D = \frac{1}{1 + (1 - y)\epsilon} \quad \text{and} \quad \eta = \epsilon \gamma y \frac{1 - (1 - y)\epsilon}{1 + \epsilon R},
\]

(20)

(21)

\[
d = D \frac{2\epsilon}{1 + \epsilon},
\]

(22)

\[
\zeta = \frac{\eta (1 + \epsilon)}{2\epsilon}.
\]

(23)

In these expressions \(\epsilon\) represents the magnitude of the virtual photon’s transverse polarisation:

\[
\epsilon = \left(1 + \frac{1}{2} \frac{y^2 + y^2\gamma^2}{1 - y - \frac{1}{2}y^2\gamma^2}\right)^{-1} \approx \frac{1 - y}{1 - y + \frac{y^2}{2}},
\]

(24)

\(D\) and \(d\) can be regarded as depolarisation factors of the virtual photon, i.e. \(D\) is the fraction of the longitudinal electron beam polarisation transferred to the virtual photon, which is emitted with a fractional energy \(y\) and at an angle \(\theta^*_\gamma\) relative to the beam direction. The quantity \(R = \sigma_L/\sigma_T\) is the
ratio of the longitudinal to the transverse virtual photon absorption cross section, which is related to the two unpolarised structure functions $F_1(x, Q^2)$ and $F_2(x, Q^2)$, measured in inclusive deep-inelastic lepton scattering from an unpolarised target, by

$$R(x, Q^2) = \frac{F_2(x, Q^2) (1 + \gamma^2) - 2xF_1(x, Q^2)}{2xF_1(x, Q^2)} .$$

(25)

From the lepton asymmetries $A_\parallel$ and $A_\perp$, the virtual photon asymmetries $A_1$ and $A_2$ can be calculated which are independent from the lepton kinematics, and are directly related to the photon–nucleon absorption cross sections for a given $x$ and $Q^2$,

$$A_1(x) = \frac{\sigma_1 - \sigma_\perp}{\sigma_1 + \sigma_\perp} = \frac{g_1(x) - \gamma g_2(x)}{F_1(x)} ,$$

(26)

$$A_2(x) = \frac{\sigma_T}{\sigma_T} = \frac{\gamma(g_1(x) + g_2(x))}{F_1(x)} .$$

(27)

Here $\sigma_1$ and $\sigma_\perp$ are the virtual photo-absorption cross sections for the projection of the total angular momentum of the photon–nucleon system along the incident photon direction $\frac{1}{2}$ or $\frac{3}{2}$, respectively. The total transverse photo-absorption cross section is given by $\sigma_T = (\sigma_1 + \sigma_\perp)/2$, while $\sigma_T$ is a term arising from the interference between transverse and longitudinal amplitudes. Since $\sigma_T \leq \sqrt{\sigma_1 \sigma_\perp}$, there is a positivity limit on the value of $A_2$:

$$A_2(x) = \frac{\sigma_T}{\sigma_T} \leq \sqrt{\frac{\sigma_T}{\sigma_T}} = \sqrt{R(x, Q^2) .}$$

(28)

Finally one can relate the spin-dependent structure function $g_1(x)$ to the longitudinal lepton asymmetry $A_\parallel(x)$ by

$$g_1(x) = \frac{F_1(x)}{1 + \gamma^2} \left[ A_\parallel(x) + (\eta - \zeta) A_2(x) \right] \approx \frac{F_1(x) A_\parallel(x)}{1 + \gamma^2} .$$

(29)

### 3.2 The spin-dependent structure functions

In the quark parton model the structure function $g_1(x)$ has a transparent probabilistic interpretation. $q^+(x)$ is defined as the quark number density for quarks of flavour $q$ in a nucleon with light cone momentum fraction $x$, and parallel (anti-parallel) orientation of the quark spin with respect to the nucleon spin. A photon with positive helicity can, due to angular momentum conservation, only be absorbed by a quark with a spin orientation anti-parallel to the photon spin, since the final state, a quark, has spin $\frac{1}{2}$ and hence cannot have spin projection $\frac{3}{2}$ (see Fig. 11). If the spin orientation of the parent nucleon is anti-parallel to the photon spin (cross section $\sigma_\perp$) one consequently probes the distribution $q^+(x)$, while in the case that the photon and the nucleon spin have the same orientation (cross section $\sigma_1$) one probes the distribution $q^-(x)$. The difference of these distributions is usually defined as spin-dependent quark distribution or quark helicity distributions, given by

$$\Delta q(x) = q^+(x) - q^-(x) .$$

(30)

$\Delta \bar{q}(x)$ is the corresponding spin-dependent anti-quark distribution.

To leading order (LO) in QCD, $g_1(x)$ can be written as a linear combination of $\Delta q(x)$ and $\Delta \bar{q}(x)$

$$g_1(x) = \frac{1}{2} \sum_q e_q^2 \left[ \Delta q(x) + \Delta \bar{q}(x) \right] ,$$

(31)

where the sum runs over the various quark flavours ($q = u, d, s, \ldots$) and $e_q$ is the corresponding quark charge in units of the elementary charge $|e|$. 
Figure 11: Interaction of a polarised virtual photon with a longitudinally polarised proton probing the quark distributions $q^- (x)$ and $q^+ (x)$.

In the unpolarised case one sums over both spin projections and obtains the ordinary quark distributions, $q(x) = q^+(x) + q^-(x)$, and the spin-independent structure functions

$$ F_1(x) = \frac{1}{2} \sum_q e_q^2 (q(x) + \bar{q}(x)) , \quad (32) $$

$$ F_2(x) = \sum_q e_q^2 x (q(x) + \bar{q}(x)) . \quad (33) $$

Furthermore, Eq. (31) can be decomposed into a flavour singlet (S) and a non-singlet (NS) component

$$ g_1(x) = g_1^S(x) + g_1^{NS}(x) , \quad (34) $$

where

$$ g_1^S(x) = \frac{1}{2} \langle e_q^2 \rangle \sum_q [\Delta q(x) + \Delta \bar{q}(x)] \equiv \frac{1}{2} \langle e_q^2 \rangle \Delta \Sigma(x) , \quad (35) $$

$$ g_1^{NS}(x) = \frac{1}{2} \sum_q e_q^2 \Delta q^{NS}(x) , \quad (36) $$

$$ \Delta q^{NS}(x) = \Delta q(x) + \Delta \bar{q}(x) - \frac{1}{N_f} \Delta \Sigma(x) , \quad (37) $$

with $\langle e_q^2 \rangle = \frac{1}{N_f} \sum_q e_q^2$ being the mean of the squared quark charges and $N_f$ the number of quark flavours taken into account.

The singlet distribution $\Delta \Sigma(x)$ is the sum of the spin-dependent quark and anti-quark distributions. This quantity is of particular importance as its first moment can be interpreted as the fraction of the nucleon’s spin originating from the spins of the quarks.

The second spin-dependent structure function $g_2(x, Q^2)$ does not have an equally transparent probabilistic interpretation. Its knowledge is required for an unambiguous determination of $g_1(x, Q^2)$ from the cross section expression given by Eq. (16). In addition, it contains further important information. From the operator product expansion [33, 34] one obtains:

$$ g_2(x, Q^2) = g_1(x, Q^2) + \int_z^1 \frac{dz}{z} g_1(z, Q^2) + \tilde{g}_2(z, Q^2) $$

$$ = g_2^{WW}(x, Q^2) + \tilde{g}_2(x, Q^2) , \quad (39) $$
where the Wandzura–Wilczek term $g_2^{WW}(x, Q^2)$ corresponds to the twist–2 part of $g_2(x, Q^2)$ and $g_2(x, Q^2)$ is a twist–3 contribution which arises only for massive quarks and is sensitive to quark–gluon correlations.

### 3.3 Moments of spin-dependent quark distributions and structure functions

Important information about the spin structure of the nucleon can be obtained from the first moments of the structure functions $g_1$ and $g_2$ similar to the unpolarised case where the low value of the first moment of $F_2$ demonstrated that quarks carry only about 50% of the nucleon’s momentum.

The first moment of the spin-dependent quark distribution is defined as

$$ \Delta q = \int_0^1 \Delta q(x) dx. \quad (40) $$

This first moment represents the fraction of the nucleon’s spin which can be attributed to the spin of the quarks of flavour $q$. The first moment of the structure function $g_1(x)$ reads in the LO-QCD parton model:

$$ I_1^p = \int_0^1 g_1(x) dx = \sum_q e_q^2 (\Delta u_q + \Delta d_q).\quad (41) $$

For the proton this yields:

$$ I_1^p = \frac{1}{2} \left\{ \frac{4}{9} (\Delta u + \Delta \bar{u}) + \frac{1}{9} (\Delta d + \Delta \bar{d}) + \frac{1}{9} (\Delta s + \Delta \bar{s}) \right\} = \frac{1}{12} \left( \alpha_3 + \frac{1}{3} \alpha_8 + \frac{4}{3} \alpha_0 \right), \quad (42) $$

with

$$ \alpha_3 = \Delta u + \Delta \bar{u} - \Delta d - \Delta \bar{d}, \quad \alpha_8 = \Delta u + \Delta \bar{u} + \Delta d + \Delta \bar{d} - 2 (\Delta s + \Delta \bar{s}), \quad \alpha_0 = \Delta u + \Delta d + \Delta d + \Delta s + \Delta \bar{s}. \quad (43) $$

Here $\alpha_3, \alpha_8, \alpha_0$ are the axial charges, i.e. the proton expectation values of the axial vector current

$$ A_{jA}^\mu = \bar{\Psi} \gamma^\mu \gamma^5 \left( \frac{\lambda_j}{2} \right) \Psi, \quad j = 0, \ldots, 8 \quad (44) $$

where the matrices $\lambda_j$ are the generators of flavour SU(3) in the Gell-Mann standard notation.

The axial charges $\alpha_3$ and $\alpha_8$ are $Q^2$ independent non-singlet quantities. They can also be obtained from the weak decay of hyperons in the spin-$\frac{1}{2}$ octet. Assuming SU(3) flavour symmetry all axial currents in this octet are given by two decay constants, $F$ and $D$. These were determined experimentally. For details see Ref. [35] and references therein. The relations are:

$$ \alpha_3 = F + D = g_A/g_V = 1.2573 \pm 0.0028 \quad \text{and} \quad \alpha_8 = 3F - D = 0.579 \pm 0.025, \quad (45) $$

where $g_A/g_V$ is the ratio of axial-vector and vector coupling constants which can be obtained from the Gamov–Teller $\beta$-decay of the neutron. The axial singlet matrix element $\alpha_0$ can, however, not be fixed from hyperon decays. Measurements of the spin-dependent structure function $g_1$ thus provide the third input for the determination of $\alpha_0$.

Using isospin symmetry ($\Delta u^p = \Delta d^p = \Delta d, \Delta d^n = \Delta u^p = \Delta u$) one obtains the corresponding expression for $\Gamma_1^p$ (of the neutron) which has a minus sign in front of the term $\alpha_3$.

The difference between the first moments of the proton and the neutron yields the famous and fundamental Bjorken sum rule [36, 37]

$$ \Gamma_1^p - \Gamma_1^n = \frac{1}{6} \frac{g_A}{g_V} = 0.209. \quad (46) $$
If one assumes that the strange quarks are not polarised, $\Delta s + \Delta s = 0$, which is equivalent to $u_0 = a_s$, one obtains the Ellis-Jaffe 'sum rule' [38]

$$\Gamma_1^{p,n} = \frac{1}{12} (F + D) + \frac{5}{36} (3F - D)$$

which yields, without QCD corrections, a value of $0.185 \pm 0.004$ for the proton and a very small value of $-0.024 \pm 0.004$ for the neutron, and a value for $\Delta^2$ of around 0.58.

Measurements on $\Gamma_1^p$, however, yielded a surprise. In the mid-eighties the value for $\Gamma_1^p$ was determined by the EMC muon experiment at CERN [39, 40] from a polarised NH$_3$ target. This experiment extended the $x$ range covered by the first spin-dependent deep-inelastic scattering experiments at SLAC [41, 42] by an order of magnitude from $x \approx 0.1$ down to $x \approx 0.01$. The value of $\Gamma_1^p$ turned out to be much smaller than the theoretical expectation, $\Gamma_1^p = 0.126 \pm 0.010$ (stat.) $\pm 0.015$ (syst.). Using the relations above including some first order QCD corrections it was concluded that the total fraction of the nucleon's spin originating from quark spins was much smaller than the expectation ($\approx 0.58$) from the relativistic quark model, i.e. $\Delta \Sigma = 0.120 \pm 0.094 \pm 0.138$. Moreover the strange quarks appeared to be polarised opposite to the parent nucleon's spin direction: $\Delta s + \Delta s = -0.190 \pm 0.032 \pm 0.046$.

### 3.4 Next-to-leading order QCD corrections to $g_1(x)$

In next-to-leading order (NLO) QCD the expression for the structure function $g_1(x)$ gets $Q^2$ dependent due to the elementary processes of gluon emission (absorption) by quarks, the creation (annihilation) of pairs of quarks and anti-quarks from (into) gluons, the coupling of three or four gluons at one vertex and higher order processes.

The expression for $g_1(x, Q^2)$ involves then the singlet quark distribution $\Delta \Sigma(x, Q^2)$, the non-singlet quark distribution $\Delta q^{NS}(x, Q^2)$ and in addition the spin-dependent gluon distribution $\Delta g(x, Q^2)$. Instead of Eq. (34) we obtain:

$$g_1(x, Q^2) = \frac{1}{2} \sum_q e_q^2 \int_x^1 \frac{dy}{y} \left[ \Delta q^{NS}(x, Q^2) \Delta C_q^{NS}(y, Q^2) + \frac{1}{N_f} \Delta \Sigma(x, Q^2) \Delta C_q^S(y, Q^2) + \Delta g \left( \frac{x}{y}, Q^2 \right) \Delta C_g(y, Q^2) \right]$$

The various coefficient functions $\Delta C_i(x, Q^2)$ can be expanded in powers of the strong coupling constant $\alpha_s(Q^2)$:

$$\Delta C_i(x, Q^2) = \Delta C_i^{(0)}(x) + \frac{\alpha_s(Q^2)}{2\pi} \cdot \Delta C_i^{(1)}(x) + \left( \frac{\alpha_s(Q^2)}{2\pi} \right)^2 \cdot \Delta C_i^{(2)}(x) + \cdots$$

At leading order $\Delta C_g^{(0)}(x) = 0$, such that the spin-dependent gluon distribution does not contribute to $g_1(x, Q^2)$. In NLO the factorisation between the helicity distributions and the coefficient functions cannot be defined unambiguously and the distributions depend on the factorisation and also the regularisation scheme. The structure function $g_1(x, Q^2)$ defined as a physical observable is of course independent of the choice of the scheme and there are straightforward transformations that relate the schemes and their results to each other. Due to the Adler-Bell-Jackiw axial anomaly [43, 44] the factorisation schemes cannot at the same time fulfil gauge invariance and conserve chiral symmetry [45]. In the gauge invariant so-called Modified-Minimal-Subtraction (MS) scheme also the first moment of the second term in the expansion (49) of $\Delta C_g(x, Q^2)$ vanishes and so $\Delta g$ does not contribute directly to the first moment $\Gamma_1$ of $g_1$. Since chiral symmetry is violated in this scheme, the quasi-quark pairs originating from gluons have the same helicity. Hence a positive gluon polarisation results in a negative sea-quark polarisation.

In the chirality conserving Adler-Bardeen [46] (AB) scheme the axial anomaly causes the first moment of $\Delta C_g^{(1)}(x)$ to be non-zero [47, 48, 49, 50, 51]. As a consequence $\Gamma_1(Q^2)$ depends on $\Delta g(Q^2)$ and...
we get:

\[ \Delta g(Q^2)_{\overline{MS}} = \Delta g(Q^2)_{\text{AB}}, \quad \Delta \Sigma(Q^2)_{\overline{MS}} = \Delta \Sigma_{\text{AB}} - N_f \frac{\alpha_s(Q^2)}{2\pi} \Delta g(Q^2). \]  

(50)

The polarised coefficient functions \( \Delta C_i(x, Q^2) \) have been calculated in the \( \overline{\text{MS}} \) scheme up to order \( \mathcal{O}(\alpha_s^2) \) \([52, 53, 54]\), the moments of the coefficient functions were calculated up to third order for the non-singlet case \([55]\) and up to second order for the singlet case \([56]\). Estimates exist for the fourth order non-singlet and the third order singlet term \([57]\). The \( Q^2 \) dependence of \( \Delta g_{\text{NS}}(x, Q^2), \Delta \Sigma(x, Q^2) \), and \( \Delta g(x, Q^2) \) can then be described (and calculated) in the selected factorisation scheme by the polarised version of the coupled Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) equations \([58, 59, 60, 61]\)

\[ \frac{d}{d \ln Q^2} \Delta g_{\text{NS}}(x, Q^2) = \frac{\alpha_s(Q^2)}{2\pi} \gamma_{\text{qg}}(Q^2) \otimes \Delta g_{\text{NS}}, \]  

(51)

\[ \frac{d}{d \ln Q^2} \left( \Delta \Sigma(x, Q^2) \right) = \frac{\alpha_s(Q^2)}{2\pi} \left( \gamma_{\text{gg}}(Q^2) + 2N_f \gamma_{\text{gg}} \right) \otimes \left( \Delta \Sigma \right), \]  

(52)

with the notation \( (a \otimes b) = \int_0^1 a(x/y) b(y) dy/y. \)

The splitting functions \( \gamma_{ij}(\frac{1}{x}) \) are in leading order QCD a measure for the probability that a parton of type \( i \) with momentum fraction \( x \) originates through one of the above elementary processes from a parton of type \( j \) with momentum fraction \( y \). They can be expanded in a power series analogue to Eq. (49) and have been calculated in NLO \([53, 62]\).

If and only if the \( x \) and \( Q^2 \) dependences of the polarised structure function \( g_1(x, Q^2) \) are experimentally determined, these relations allow in principle to determine the shape of the individual spin-dependent parton distributions. This is discussed in some detail in Sect. 4.2.
4 Experimental determination of the spin-dependent structure function \( g_1 \) and of the spin-dependent quark distributions

4.1 The spin-dependent proton and neutron structure functions \( g_1^{p,n}(x) \)

After the observation by EMC [39, 40] that only a small fraction of the nucleon spin can be attributed to the spins of the quarks the spin-dependent structure functions \( g_1 \) and \( g_2 \) have been investigated in much detail by several experiments at SLAC, the SMC muon experiment at CERN and the HERMES experiment at DESY. All these experiments [52, 53, 54] were only possible due to enormous technological achievements in the development of polarised targets and polarised beams during the last two decades.

A comparison of the various experiments and the relevant beam and target parameters can be found for example in recent reviews [63, 64, 31, 65]. As is shown below these measurements confirm the original EMC finding but with now much better statistical and systematic accuracy.

For most of the beam time available at HERMES in the years 1995-2000 longitudinally polarised hydrogen, deuterium and \(^3\)He atomic gas targets have been used. Especially the pure atomic hydrogen and deuterium targets have the big advantage of high nuclear polarisation of about \( P_T \approx 0.9 \) not diluted by unpolarisable nuclei encountered in solid polarised targets. The experimental asymmetry \( A_{\text{exp}} \) is, if the small contribution of \( A_2 \) to \( A_1 \) is neglected, related to the virtual photon asymmetry \( A_1 \) by

\[
A_{\text{exp}} = f \cdot P_T \cdot D \cdot P_B \cdot A_1,
\]

and the statistical uncertainty for the asymmetry is given by

\[
\delta A_1 = \frac{1}{\sqrt{N^+ + N^-}} (f \cdot P_T \cdot D \cdot P_B)^{-1},
\]

where \( N^+ (N^-) \) is the counting rate for the target spin parallel (anti-parallel) to the beam spin, ~\( f \) is a dilution factor describing the fraction of deep-inelastic events originating from polarised nucleons in the target, \( P_T \) and \( P_B \) are the target and beam polarisations, and \( D \) is the depolarisation factor (Eq. (20) in sect. 3.1), i.e. the fraction of the longitudinal electron beam polarisation transferred to the virtual photon. Note that for typical values of these quantities for the EMC experiment (\( f \approx 0.15, P_T \approx P_B \approx 0.8, D \approx 0.5 \)) \( \delta A_1 \) is about 20 times bigger than the corresponding uncertainty from the counting rate alone. It is obvious that the best way to decrease this uncertainty substantially is to increase \( f \). For atomic hydrogen and deuterium targets \( f = 1 \), compared to \( f = 3/17 \) for ammonia, NH\(_3\), or \( f \approx 6/20 \) for deuterated ammonia, ND\(_3\), targets.

Generally one extracts the cross section asymmetry \( A_1 \) from the measured counting rates \( N^+ \) and \( N^- \), defined above, using the relation

\[
\frac{A_1}{D} = \frac{N^+ L^+ - N^- L^-}{D(N^- L_P^+ + N^+ L_P^-)} ,
\]

where \( L^\pm \) are the dead-time-corrected luminosities for each target spin state, and \( L_P^\pm \) are the dead-time-corrected luminosities weighted by the product of the magnitudes of the beam and target polarisations for each spin state. The longitudinal virtual photon asymmetry \( A_1 \) is determined from \( A_1 \) using Eq. (18) in section 3.1. For the transverse asymmetry \( A_2 \), entering this equation, different assumptions have been made which in the course of the years were refined with increasing experimental knowledge about its kinematical behaviour.

Finally one determines the spin structure function \( g_1(x, Q^2) \) from the asymmetries \( A_1 \), or \( g_1/F_1 \) respectively, using the prescriptions given in section 3.1 and parameterisations of the structure function \( F_2(x, Q^2) \) [66] and the ratio \( R(x, Q^2) = \sigma_L/\sigma_T \) based on high statistics unpolarised data [67] and the procedure for radiative corrections from Ref. [68].
4.1.1 $g_1^p$ from $^3\text{He}$

In 1995 data taking started at HERMES using the polarised $^3\text{He}$ target. The average value of the target polarisation during the experiment was 46 % with a fractional uncertainty of 5 %. The target polarisation was reversed every 10 minutes by reversing the helicity of the laser light. The average beam polarisation for the analysed data was 55 % with a fractional systematic error of 5.4 %. After applying data quality criteria and kinematic cuts ($E^2 > 1 \text{ GeV}^2$, $W^2 > 4 \text{ GeV}^2$ and $y > 0.85$) 2.7 · $10^6$ events out of the collected 4.2 · $10^6$ events were available for the asymmetry analysis.

The nuclear wave-function of $^3\text{He}$ is dominated by the configuration with the protons paired to zero spin. Therefore, most of the $^3\text{He}$ asymmetry is due to the neutron [69] and the $^3\text{He}$ target can be considered to a good approximation as a polarised neutron target. A correction for the non-zero polarisation of the protons (-0.028 ± 0.004) and the neutron polarisation (0.86 ± 0.02) [70] has to be applied to the $^3\text{He}$ results to obtain those for the neutron. The virtual photon asymmetry $A_1^p$ has been determined from $A_1$ using Eq. (18) with the assumption $A_2 = 0$.

![Figure 12](image)

Figure 12: The polarised virtual photon asymmetry $A_1^p$ (upper panel) measured with longitudinally polarised $^3\text{He}$ targets and the polarised structure function $g_1^p$ (lower panel) as a function of $x$.

The HERMES results for $A_1^p(x)$ and $g_1^p(x)$ [71] are shown in Fig. 12 together with those from the SLAC experiments E142 [72] and E154 [73]. The agreement between all three experiments is excellent, proving that the systematic uncertainties of the different experiments are well under control. The E154 data extend to lower values of $x$ due to their higher beam energy of 48.3 GeV. They have much smaller error bars than those from HERMES and E142. The asymmetry and the spin-dependent structure function are negative over essentially the whole $x$ range covered by the data. At $x \approx 0.5$ they are compatible with zero or slightly positive. Recent preliminary high statistics data from TJNAF [74], which cover the range $0.3 < x < 0.6$, indicate that the asymmetry is positive and rising for $x > 0.5$, but the behaviour in the large $x$ range, $x > 0.6$, is essentially unknown. The behaviour in the measured $x$ range can be interpreted using the QPM relation for $g_1(x)$ (see Eq. (31)). It tells us, that the positive spin-dependent down-quark distribution in the neutron ($\Delta d^n(x) \equiv \Delta u(x)$) is largely cancelled by the much smaller negative spin-dependent up-quark distribution in the neutron ($\Delta u^n(x) \equiv \Delta d(x)$) which
enters the cross section with a four times larger charge weight, or even more specific: $4\Delta d(x) < -\Delta u(x)$ in the measured $x$ range.

4.1.2 $g_1^p$ from hydrogen

In 1996 and 1997 data were taken at HERMES with the longitudinally polarised atomic hydrogen target. After all kinematic and geometrical cuts were applied and the background subtracted about 2.5 million DIS events were available for the determination of $g_1^p$. Detailed studies of the target polarisation and the fraction of molecules, generated by recombination on the walls of the storage cell, have been performed as a function of several parameters like storage cell temperature and static holding field and compared to extensive Monte Carlo simulations and rate equation calculations [14, 15, 16, 17]. All these studies resulted in a very good understanding of the target polarisation. The spin direction in the target can be reversed in less than a second by selecting different spin states in the ABS. In operation the length of the time reversals between spin states was randomised and was in the order of one minute. The mean absolute polarisation value used for the analysis was $\langle |P^T| \rangle = 0.868$ with an overall systematic uncertainty of 5.5%. The average beam polarisation value of the analysed data was $\langle |P^B| \rangle = 0.55 \pm 0.02 \, (\text{sys.})$, where the systematic error is dominated by a scale uncertainty obtained from polarisation rise-time calibrations. The distribution of the events in the kinematic $Q^2 - \nu$ plane is shown in the left panel of Fig. 13.

For the published results [75] the kinematic requirements used in the analysis were: $Q^2 > 0.8 \, \text{GeV}^2$, $W > 1.8 \, \text{GeV}$ and $0.1 < y < 0.85$, the resulting lowest $x$ value being 0.021. In a new analysis the kinematic range has been extended to lower values of $x$ ($x > 0.0021$) and $Q^2$ ($Q^2 > 0.1 \, \text{GeV}^2$), by increasing the upper limit in $y$ from 0.85 to 0.91. Fig. 13 (right panel) shows the distribution of the $x$ and $y$ bins in the kinematic $x$-$y$ plane, together with the different kinematical and geometrical boundaries.

![Figure 13](image-url)

Figure 13: Distribution of the collected DIS events in the $Q^2$, $\nu$ plane (left panel); distribution of bins used for the analysis in the kinematical $y$-$x$ plane (right panel). The $x$ bins are evenly spaced in $\ln x$. The filled circles indicate the bin centres, the curved lines indicate the kinematical requirements $40 \, \text{mrad} \leq \theta \leq 220 \, \text{mrad}$ and $W > 1.8 \, \text{GeV}$.

At the high $y$ values of the low-$x$ bins the momenta of the scattered electrons are correspondingly small and therefore such an analysis requires good understanding of the electron detection at low momenta, especially the trigger efficiency and the steeply rising contributions from background and radiative events. For these low $x$, low $Q^2$ data most of the relevant contributions to experimental systematic uncertainties originate from normalisation, corrections for background from hadrons, wrongly identified as positrons, and from charge symmetric processes, acceptance cut variations and scattering angle systematics. The systematic uncertainties arising from radiative and smearing corrections have
been minimised using an iterative procedure for the asymmetry at the Born cross section level. In the most recent analysis the application of the radiative corrections was revised, which mainly enlarged the statistical errors at low $x$ due to the proper subtraction of the radiative background contribution.

The proton structure function ratio $g_1/F_1$ has been determined from the measured count rates using Eqs. (55) and (29). While in the case of the neutron data $A_2^p$ was determined with the assumption $A_2^p = 0$, the contribution of the transverse virtual photon absorption asymmetry $A_2^p$ was taken into account using a fit to existing data from SLAC [76, 77] and SMC [78, 79] of the form $A_2^p = 0.5 \times x/\sqrt{Q^2}$. The published results for $x > 0.021$ and the new preliminary data points for the lower $x$ region, $0.0021 < x < 0.021$ are shown in Fig. 14 as a function of $x$, together with the corresponding data from SMC [80, 81], SLAC-E143 [77] and SLAC-E155 [82].

![Proton HERMES preliminary](image)

Figure 14: HERMES results for the structure function asymmetry $g_1/F_1$ (left panel) and the structure function $g_1^p$ (right panel) as a function of $x$, together with the corresponding data from E-143 and SMC in the $x$ range covered by the HERMES data. The HERMES data points for $x < 0.021$ are preliminary.

The statistical error bars are small as are the systematic errors shown by the error bands. It should be noted that the final statistical errors for the preliminary data points will be substantially larger (increasing with decreasing $x$ from about a factor of 1.5 to about a factor of 2.5) than the shown error bars due to a revised treatment of radiative corrections.

For each point in $x$ the mean $Q^2$ value is different for each experiment as can be seen from the lower panel of the figure. Due to the higher beam energy the SMC data have, in every $x$ bin, typically about an order of magnitude higher ($Q^2$) values than the HERMES and E143 data. The SMC small $x$ ($Q^2 > 1$ GeV$^2$) data are confirmed for the first time. No significant $Q^2$ dependence of the structure function ratio $g_1/F_1$ (and also the virtual photon asymmetry $A_1$) within the experimental uncertainties can be seen [83].

Therefore the $Q^2$ dependence of $g_1$ and $F_1$ must be very similar over the whole $x$ range and one must conclude that even at very low $Q^2$ values, where the applicability of pQCD is questionable, QCD radiative corrections and higher-twist contributions must have similar relative magnitude in $g_1$ and $F_1$, such that they cancel in the ratio.
The asymmetry decreases continuously to small positive values close to zero at very small values of $x$. On the other hand, it rises smoothly to a value of $A \approx 0.8$ at the highest $x$ point ($x \approx 0.8$) in agreement with the expectation $A_1 \to 1$ for $x \to 1$ (i.e. the quark carrying most of the nucleon's momentum is also carrying most of its spin). The $x$ dependence of $g_1^d(x)$ at the measured values of $Q^2$ is shown in the right panel of Fig. 14. At $x < 0.1$ the spin-dependent structure function is essentially constant with a value of about 0.3 at these low values of $Q^2$. The difference between the HERMES and the SMC data is due to the above discussed $Q^2$ dependence of $g_1(x, Q^2)$.

### 4.1.3 $g_1^d$ from deuterium

In the years 1998 to 2000 data were taken at HERMES with the longitudinally polarised deuterium target. Altogether about 9 million DIS events were collected, corresponding to an integrated luminosity of about 80 pb$^{-1}$. The average beam polarisation was $\langle |P^B| \rangle = 0.53$ and the average target polarisation $\langle |P^T| \rangle = 0.845$. As in the proton case, discussed above, the kinematic range of the deuteron data is $0.0021 < x < 0.85$, $0.1 < y < 0.91$, $Q^2 > 0.1$ GeV$^2$ and $W^2 > 3.24$ GeV$^2$.

Since the available deuteron data for $g_1^d(x)$, including the most recent ones from SLAC experiment E155x [84], do not indicate a significant deviation from the Wandzura-Wilczek twist-2 expression [33] (see Eq. (39) in Sect. 3.1), $A_2^d$ was calculated using this relation and parameterisations for $g_1/F_1$, $F_2$ and $R$:

$$A_2^{WW} = \frac{\gamma x}{1 + \gamma^2} \frac{1 + R(x, Q^2)}{F_2(x, Q^2)} \int_x^1 dz \frac{1 + \gamma^2}{z^2} \frac{F_2(z, Q^2)}{1 + R(z, Q^2)} \frac{g_1^d(z, Q^2)}{F_1^d(z, Q^2)}$$

(56)

The difference between the corresponding values and those obtained for the assumption $g_1^d = 0$, which corresponds to $A_2 = \gamma \cdot g_1^d/F_1^d$, was used as the contribution due to $A_2$ to the systematic uncertainty.

Figure 15: Preliminary HERMES results for the structure function asymmetry $g_1^d/F_1^d$ as a function of $x$ (left panel) together with the corresponding data from E143, E155 and SMC in the $x$ range covered by the HERMES data. The right panel shows the preliminary HERMES results for the structure function $g_1$ together with the SMC results at the measured values of $Q^2$ as a function of $x$ in the $x$ range covered by the HERMES data.
The preliminary HERMES result for the structure function ratio $g_1^u/F_1^u$ for the data from the year 2000 is shown in the left panel of Fig. 15 together with the corresponding data from SMC [80], E143 [77] and E155 [85]. The seven data points from SMC at even lower values of $x$ down to $x = 10^{-4}$ are not shown here. No corrections for smearing effects have been applied yet (but the revised treatment of radiative corrections was already implemented) to these preliminary HERMES data. They have substantially smaller error bars, show much smaller point to point fluctuations than those from the other experiments and determine the $x$ dependence of this ratio now very well. This is specially evident in the range of very small $x$, where the statistical quality of the data has improved by nearly an order of magnitude, but also at higher $x$ where the central values of various previous data sets differed substantially.

At very small $x$ the asymmetry is compatible with zero, with a slight tendency to negative values. One can conclude from this behaviour that at these small $x$ values $\Delta d(x)$ and $\Delta u(x)$ have about the same magnitude, but opposite sign. If the asymmetry really gets negative then one can speculate that either the absolute value of the spin-dependent down quark distribution gets bigger than that of the spin-dependent up quark distribution or that the net spin-dependent sea quark distribution is negative. At large $x$ the rise of the asymmetry with $x$ is much less steep than for the proton data, the limit at $x \to 1$ possibly being substantially smaller than unity. As in the proton case there is no significant $Q^2$ dependence of the asymmetry within the experimental errors.

---

**Figure 16.** $Q^2$ dependence of $g_1(x, Q^2)$ for $Q^2 > 1$ GeV$^2$ (left panel). To evaluate the $Q^2$ dependence, the data have been shifted to common $x$ values. The right panel shows the corresponding $Q^2$ dependence of the unpolarised structure function $F_2(x, Q^2)$ in a similar $x$ range.

The preliminary HERMES result from the 2000 data for $g_1^u(x)$ at the measured values of $Q^2$ is
shown as a function of $z$ together with the corresponding data from SMC in the right panel of Fig. 15. It is evident that with these new data also the $x$ dependence of $g_1^p$ is very well determined in the range $x > 2 \cdot 10^{-2}$. At lower $x$ it is compatible with zero with a slight tendency to negative values. Since the ratio $g_1/F_1$ appears to be independent of $Q^2$ while the structure function $F_1$ (or $F_2$ and $R$ equivalently) shows significant scaling violations, $g_1$ is also expected to vary with $Q^2$ correspondingly. This is demonstrated for the proton case in Fig. 16 (left panel) where the $Q^2$ dependence of $g_1(x, Q^2)$ for fixed values of $x$ is shown for all available data with $Q^2 > 1 \text{ GeV}^2$. For this comparison the data of the various experiments have been shifted to common $x$ values. The $Q^2$ dependence of $g_1$ is fairly well described by a phenomenological fit of $g_1/F_1$ [83]:

$$
g_1/F_1 = \frac{x^{0.7}(0.817 + 1.014x - 1.489x^2)(1 - 0.04/Q^2)}{F_1^{\text{P}}} [86] \text{ and } R [67], \text{ but also by a NLO-QCD fit. For comparison the } Q^2 \text{ dependence of the high precision}
$$

![Figure 17](image_url)

Figure 17: Compilation of recent data on the spin-dependent structure functions $xg_1(x, Q^2)$ for proton, deuteron and neutron ($^3\text{He}$). All data are given at their quoted mean $Q^2$ values. Also shown are very recent preliminary results for $xg_1^n$ at large $x$ from an experiment at TJNAF [74] (courtesy U. Stöslein).

data for the unpolarised structure function $F_2$ from fixed target and collider experiments [87, 88, 89, 90] is shown for the kinematic range $x > 0.008$ in Fig. 16 (right panel). It is evident that as compared to the unpolarised data, polarised DIS data are considerably less precise and cover a much smaller $Q^2$ range. Nevertheless some conclusions can be drawn from the left panel of Fig. 16. The scaling violations agree within errors with the expectations from pQCD, they are weak for $x > 0.15$ but significant and positive for $x < 0.15$.

The present status of inclusive spin-dependent DIS is summarised in Fig. 17, where the quantity $xg_1$ is plotted vs $x$ for $x > 0.0001$. All the data are displayed at the measured $Q^2$ values and the plotted errors are the quadratically summed systematic and statistical uncertainties.

### 4.2 QCD analysis of $g_1(x, Q^2)$

As has been discussed in Sect. 3.4, one can in principle extract the spin-dependent parton distributions and their integrals from the $x$ and $Q^2$ variation of $g_1(x, Q^2)$. However this is only possible if the data
measured in inclusive polarised deep-inelastic scattering span a sufficiently large range in $x$ and $Q^2$ and the error bars are small enough. The experimentally measured quantity $A_1 \approx g_1/F_1$ is, however, independent of $Q^2$ within errors, and $g_1$ evolves basically like $F_1$ implying that very little additional information can be extracted from a polarised DGLAP fit. Compared to the precise measurements of the unpolarised structure function $F_2$, the corresponding data in the polarised case are relatively poor. The $x$ and $Q^2$ range covered by the present data for $g_1(x, Q^2)$ is very limited, as can be seen from Fig. 16. Each of the experiments gives essentially only a few $Q^2$ points per $x$-bin. A DGLAP fit requires an experimentally well measured $x$ dependence of $g_1$ at a fixed value $Q^2_0$ that is used as a starting point for the $Q^2$ evolution. This is not yet available and also the statistical accuracy of the present data is hardly sufficient to reliably determine the predicted variations with $\ln Q^2$. Nevertheless several authors have performed fits to the existing data on $g_1(x, Q^2)$ [91, 92, 93, 51, 80, 94, 95, 83, 96, 97, 98, 99, 100, 101, 102]. Depending on the data sets and the assumptions used the extracted parton distributions differ quite significantly but not the main conclusion. The results improved a lot over the years with increasing quality of the experimental data. As an example, the results from the most recent analysis [102] in NLO-QCD are briefly reported, which is based on all available experimental data apart from the HERMES deuterium data and the recent $g_2$ data from E155x. Fig. 18 shows the results for the fitted spin-dependent distributions for up and down valence quarks, gluons and sea-quarks from this analysis (BB) together with the results from Ref. [92] (GRSV) and from Ref. [98] (AAC) at a $Q^2$ value of 5 GeV$^2$. The dotted lines on the right hand side of the histograms indicate the positivity bounds choosing the unpolarised distributions [103] for reference. The extracted moments have also been compared to the results of Ref. [51]. To reduce the number of fit parameters a flavour symmetric spin-dependent sea-quark distribution has been assumed and the first moments of the valence spin distributions $\Delta u_v$ and $\Delta d_v$ have been fixed by using the $F$ and $D$ constants experimentally determined from weak-decay of the members of the spin-$\frac{1}{2}$ baryon octet. This is a very strong constraint. Whether it is really justified has to be investigated experimentally with semi-inclusive data, which is discussed below.

On the basis of these fits the spin-dependent valence distributions appear to be rather well determined, especially the up quark valence distribution, which dominates the measured asymmetries due to

Figure 18: Spin-dependent parton distributions at $Q^2 = 5$ GeV$^2$ from updated NLO-QCD fits in the MS scheme (from Ref. [102]).
the fact that $e^2_v = 4e^2_d = 4e^2_s$. The distribution $\Delta u_v$ is positive over the whole $x$ range and still non-zero at $x = 10^{-3}$. The fits of the three groups agree very well. The error bars are substantially larger in the case of the down valence quark, as can be seen from the shaded area representing the full correlated error bands of the fit. The $\Delta d_v(x)$ distribution is negative over the whole $x$ range, the absolute value in the maximum of the $\Delta d_v$ distribution is about three times smaller than the corresponding value for $\Delta u_v(x)$. A very small negative sea-quark polarisation is obtained. The $x$-weighted distribution $x\Delta q(x)$ has a maximum absolute value of about 0.01 at an $x$ value around 0.1, which makes it very difficult to be determined experimentally. The biggest differences between the various fits and also the biggest uncertainties occur for the spin-dependent gluon distribution. The distribution $x\Delta g(x)$ is positive over the indicated $x$ range at $Q^2 = 5$ GeV$^2$, but its magnitude is uncertain by nearly an order of magnitude. The authors conclude from their fits that the present data do not contain significant higher twist contributions in the range $Q^2 > 1$ GeV$^2$, similar to an analysis of E155 [83].

The moments of the spin-dependent parton distributions extracted from the various fits agree rather well. They suffer, however, substantially from uncertainties due to the extrapolation to small $x$. The moment $\Delta u$ is in the range 0.73 to 0.86 (with an error of about 10%) substantially smaller than the value $\frac{2}{3}$ expected from the naive QPM, and the moment $\Delta d$ ranges from $-0.40$ to $-0.46$ (with an error of 25-30 %), somewhat larger in absolute value than the QPM expectation of $-\frac{1}{3}$. All authors obtain a negative small sea-polarisation with a value for $\Delta q_s$ between $-0.04$ and $-0.09$ for each quark and anti-quark flavour.

Consequently these fits result in a relatively small value of $\Delta \Sigma$

$$\Delta \Sigma \approx 0.14 \ldots 0.20,$$

which is very similar to the original finding of EMC. A substantial positive moment of the spin-dependent gluon distribution is also found:

$$\Delta g \approx 0.68 \ldots 1.26.$$ (58)

In this quantity the errors of the fits are large, up to about 70 %, and also the results of the various groups differ most. As stated above, the polarised gluon distribution is essentially unconstrained by the present inclusive data on polarised deep-inelastic scattering.

If one takes the values of the moments on face value and compare them to the angular momentum sum rule for the proton spin

$$\frac{1}{2} = \frac{1}{2} \Delta \Sigma + L_z^q + L_z^g,$$ (59)

one sees that the sum of the spin contributions of quarks and gluons is much larger than the expectation value of the proton spin. Consequently the orbital angular momentum contributions $L_z = L_z^q + L_z^g$ must be substantial and negative

$$L_z \approx -0.3 \ldots -0.9.$$ (60)

Note that already nearly 30 years ago it has been demonstrated that orbital angular momentum contributions of quarks could explain [104] the deviation of the measured value of $g_A/g_V$ from the quark model prediction $5/3$. An experimental determination of the orbital momentum contribution is very challenging. At present one possibility being discussed is an indirect determination via moments of generalised parton distributions. This is discussed in Sect. 5.1.

4.3 The generalised Gerasimov-Drell-Hearn integral

The Gerasimov-Drell-Hearn (GDH) sum rule [105, 106] relates the anomalous contribution $\kappa$ to the magnetic moment of the nucleon with the total absorption cross sections for circularly polarised real photons on polarised nucleons. It is written as:

$$\int_{v_0}^{\infty} \left[ \sigma_{1/2}(\nu) - \sigma_{3/2}(\nu) \right] \frac{dv}{v} = -\frac{2\pi^2 \alpha}{M^2} \kappa^2,$$ (61)

where $\sigma_{1/2(3/2)}$ is the photo-absorption cross section for the total helicity of the photon-nucleon system equal to 1/2 (3/2) and $v_0$ is the threshold for the lowest lying inelastic channel. The GDH sum rule
holds without any assumption on the type of target, i.e. it is equally valid for protons, neutrons and nuclei. It is derived from the Compton forward-scattering amplitude following the general principles of Lorentz and gauge invariance while it is non-perturbative in nature. For the proton ($\kappa_p = +1.79$) the GDH sum rule prediction is $-204$ $\mu$b, for the neutron ($\kappa_n = -1.91$) it is $-233$ $\mu$b, and for the deuteron ($\kappa_d = -0.143$) it is $-0.65$ $\mu$b.

The generalisation of the GDH integral to virtual photons with energy $\nu$ and squared four-momentum $-Q^2$ connects the static ground state properties of the nucleon with its helicity structure as measured in the resonance and deep-inelastic regions. It reads:

\[
I_{\text{GDH}}(Q^2) = \int_{0}^{\infty} \left[ \sigma_{1/2}(\nu, Q^2) - \sigma_{3/2}(\nu, Q^2) \right] \frac{d\nu}{\nu}
= \frac{8\pi^2\alpha}{M} \int_{0}^{\infty} \frac{g_1(x, Q^2) - \gamma^2 g_2(x, Q^2) dx}{K}
= \frac{8\pi^2\alpha}{M} \int_{0}^{\infty} \frac{A_1(x, Q^2) F_1(x, Q^2) dx}{K},
\]

where $K$ is the virtual-photon flux factor, which in the Gilman convention [107] is $K = \nu \sqrt{1 + \gamma^2}$. Examining the $Q^2$ dependence of the GDH integral provides a means to study the transition from polarised real-photon absorption ($Q^2 = 0$) on the nucleon to polarised deep-inelastic scattering.

In the deep-inelastic limit, i.e. $\gamma^2 \ll 1$, Eq. 62 simplifies to

\[
I_{\text{GDH}}(Q^2)_{\gamma^2 \to 0} = \frac{16\pi^2\alpha}{Q^2} \int_{0}^{1} g_1(x, Q^2) dx - \frac{16\pi^2\alpha}{Q^2} \Gamma_1.
\]

Since the measured value of $\Gamma_1^p$ is positive for higher $Q^2$, $I_{\text{GDH}}^p(Q^2)$ must change sign at some lower $Q^2$ in order to reach the negative value predicted by the GDH sum rule at the real-photon point $Q^2 = 0$. $\Gamma_1^n$ is negative for all measured $Q^2$.

Tests of the GDH sum rule require suitable polarised real-photon beams and polarised targets. Only recently first results from an experiment at MAMI/Mainz with a polarised proton target and photon energies below 800 MeV have been published [108]. The measurements have been extended to an energy of 3.0 GeV at ELSA/Bonn [109] and further real photon experiments are underway at SLAC [110] and TJNAF [111]. The difference of cross sections $\sigma_{3/2}$ and $\sigma_{1/2}$ shows a very pronounced resonance behaviour in the first (delta), second and third resonance region followed by a smooth Regge type energy dependence up to 3 GeV.
Several experiments to measure the generalised GDH integral at low and intermediate $Q^2$ are underway at TJNAF [112]. They cover, however, only the resonance region.

The deep-inelastic contribution to the generalised GDH integral has been investigated, for the first time, by HERMES, using the polarised $^3$He, hydrogen and deuterium data discussed in the previous section. The HERMES experiment is ideally suited for such measurements due to the large angular acceptance. As shown in Fig. 13 the deep-inelastic events are, above the minimum angle of 40 mrad, continuously distributed over the kinematic plane allowing - without interpolation - the integration over $\nu$ for several bins in $Q^2$. The results from $^3$He and H have already been published [113, 114], a publication covering also the deuterium data is in preparation. Details of the analysis can be found in these publications. Here only the main results are summarised.

![Figure 20](image)

**Figure 20:** a) $I_{\text{GDH}}$ as a function of $Q^2$ for various upper limits of integration: $W^2 = 5.42 \text{ GeV}^2$ (triangles), $W^2 = 45 \text{ GeV}^2$ (squares) and the total integral (circles). b) $I_{\text{GDH}} \cdot Q^2/(16\pi^2\alpha)$ as a function of $Q^2$.

Fig. 19 shows the generalised GDH integral for the proton and the neutron as a function of $Q^2$, determined in the kinematic range $0.8 \text{ GeV}^2 \leq Q^2 \leq 12 \text{ GeV}^2$ and $W > 1.8 \text{ GeV}$ [113]. The sizes of the systematic errors are indicated by the bands on the bottom of the figure. The integral for the proton is positive and falling with $Q^2$, the integral for the neutron is negative over the whole $Q^2$ range. The data are compared with estimates of the integrals, derived using a parameterisation of $A_1$ given by a NLO QCD analysis [115], that do not include contributions from nucleon resonances or higher-twist effects. The dashed curves show the integrals for this parameterisation over the $x$ range measured by HERMES. They are in good agreement with the data, indicating that higher-twist effects do not contribute significantly in the deep-inelastic region, even at the lowest measured $Q^2$ values. The dashed-dotted curves show the integrals of the NLO parameterisation from $x_0$ to infinity, thus including in addition the yet unmeasured low $x$ contribution. The contribution of the nucleon-resonance region has been studied for the proton with a separate data sample which covered the $W$ range $1 \text{ GeV}^2 \leq W^2 \leq 4.2 \text{ GeV}^2$ and $1.2 \text{ GeV}^2 \leq Q^2 \leq 12 \text{ GeV}^2$ [114]. After applying data quality criteria about 0.13 million events were selected.

Fig. 20a shows the contribution $I_{\text{GDH}}^{\text{res}}$ of the resonance region ($W^2 \leq 4.2 \text{ GeV}^2$) to the integral (triangles), the full measured integral $I_{\text{GDH}}^{\text{meas}}$ ($W^2 \leq 45 \text{ GeV}^2$) (squares) and the total integral $I_{\text{GDH}}$ (circles). The latter was calculated in each $Q^2$ interval by adding to $I_{\text{GDH}}^{\text{res}}$ an estimate of the unmeasured region $W^2 > 45 \text{ GeV}^2$ using a multiple-Reggeon exchange parameterisation [116] for $\sigma_{1/2}(\nu, Q^2) - \sigma_{3/2}(\nu, Q^2)$ at high energy. The curve is a model [117] for the total integral, based on a $Q^2$ evolution
of $g_1$ and $g_2$ without consideration of any explicit nucleon-resonance contribution. The contribution of the resonance region to the whole measured integral is about 50% at $Q^2 < 2$ GeV$^2$, it decreases rapidly with $Q^2$ and is small for $Q^2$ values above 3 GeV$^2$. In the whole energy range $I_{GDH}$ has a simple $1/Q^2$ dependence. This is demonstrated in Fig. 20b, where the results for $I_{GDH}$ are multiplied by $Q^2$. As can be seen from Eq. (63) this product will be constant if there are no large effects from either resonances or non-leading-twist. This behaviour is reflected by the data. Moreover, the data indicate that the sign change of $I_{GDH}$ to meet the real-photon limit occurs at $Q^2$ values lower than 1.2 GeV$^2$.

Finally Fig. 21 shows the preliminary results for the $Q^2$ dependence of the generalised GDH integral $I_{GDH}$, summed over the nucleon-resonance and the deep-inelastic region, for the proton, the deuteron and the neutron, extracted from the polarised hydrogen and the preliminary deuterium data. The results for the neutron agree with those obtained from the $^3$He data, but with much better statistical accuracy. These data will allow to study also the $Q^2$ dependence of the generalised GDH integral for the proton-neutron difference, which for large $Q^2$ is expected to obey the Bjorken sum rule.

4.4 Spin-dependent parton distributions from semi-inclusive DIS

4.4.1 Semi-inclusive asymmetries

As has been discussed in Sect. 4.2 the spin-dependent up-quark distribution $\Delta u(x)$ seems to be rather well determined from the QCD fits as well in shape as in magnitude, while the fits give a much larger possible range for the spin-dependent down quark distribution $\Delta d(x)$. The spin-dependent sea quark distributions and especially $\Delta s(x)$ are essentially undetermined and could either have a positive or a negative sign, and very different solutions are possible for the spin-dependent gluon distribution $\Delta g(x)$ from such fits. A detailed understanding of the spin structure of the nucleon requires, therefore, that the $x$-dependence and the moments of the spin-dependent quark distributions, separated into contributions from quarks and anti-quarks of the various flavours, and the spin-dependent gluon distribution are determined separately. This can in principle be achieved in semi-inclusive spin-dependent deep-inelastic scattering when in addition to the scattered lepton also a leading hadron is detected [118]. Since the charge and the identity of high energy forward hadrons are correlated to the flavour of the struck quark (see Fig. 22), this type of reaction can be used to study the flavour dependence of the spin-dependent quark distribution. Detection of hadrons provides a statistical tag on the struck quark flavour. Precise spin asymmetries allow a separation of the individual contributions of each quark flavour to the nucleon spin.

Provided factorisation of the hard scattering process and the fragmentation process holds, the cross
section for the production of a particular hadron $h$ can be written in leading order (LO) QCD as

$$\sigma^h(x, Q^2, z) \propto \sum_q e_q^2 D_q (z, Q^2) , \tag{64}$$

where the fragmentation function $D_q (z, Q^2)$, which parameterises all our ignorance of the non-perturbative fragmentation process, is the probability density that a struck quark of flavour $q$, probed at the scale $Q^2$ fragments into a hadron $h$ with energy $E_h$, or fractional virtual photon energy $z = E_h/\nu$, respectively.

The fragmentation functions are normalised to conserve energy and particle multiplicities. Their $Q^2$ dependence is well described in LO and NLO QCD [119, 120, 121]. In LO QCD and under the assumption that contributions from the structure function $g_2$ vanish, the semi-inclusive double spin asymmetry $A_1^h$ for the production of a hadron of species $h$ is related to the spin-dependent and independent quark distributions and the fragmentation functions by:

$$A_1^h (x, Q^2) = C_R (x, Q^2) \sum_q e_q^2 \Delta q (x, Q^2) \frac{\int_{x_{\min}}^1 D_q (z, Q^2) \, dz}{\sum_{q'} e_{q'}^2 \Delta q' (x, Q^2) \frac{\int_{x_{\min}}^1 D_{q'} (z, Q^2) \, dz}{}} , \tag{65}$$

where the integration is performed over the range in $z$ where one can assume that the hadrons originated from the struck quark and not from the target remnant. In Eq. (65) it is assumed that the spin-dependent and spin-independent fragmentation functions are the same. In the HERMES analysis it has been assumed that this is fulfilled for pseudo-scalar mesons like pions and kaons. There are, however, also severe concerns about this assumption [122]. The correction factor $C_R (x, Q^2) = (1 + R (x, Q^2)) / (1 + \gamma^2)$ accounts for the longitudinal component of the photon nucleon cross section that is included in experimentally determined parameterisations of $q (x, Q^2)$ but not in $\Delta q (x, Q^2)$. The inclusive asymmetry (where the 'effective fragmentation functions' are unity) can be similarly expressed by $A_1^h (x, Q^2) = C_R (x, Q^2) \sum_q e_q^2 \Delta q (x, Q^2) / \sum_{q'} e_{q'}^2 \Delta q' (x, Q^2)$.

A crucial question is of course whether the observed hadron really contains information about the struck quark, or whether in reality it originated from one of the target remnants. Based on the study of correlations between hadrons from the current and the target fragmentation regions by EMC [123] it has been argued [124] that at least four units of rapidity

$$\eta = \frac{1}{2} \ln \frac{E^h + p_t^h}{E^h - p_t^h} \tag{66}$$

where $p_t^h$ is the hadron momentum along the beam direction, are needed for a clean separation of these regions (Berger’s criterion [125]). This would mean that in general the HERMES centre-of-mass energy
of about 7 GeV is too low for such a separation. In principle, this issue is indeed a question of concern. In practice, however, unpolarised semi-inclusive data from HERMES demonstrate that the problem is not so important, as shown by the measurements discussed in the next two subsections.

### 4.4.2 Pion multiplicity distributions

In Fig. 23 (left panel), the measured $z$ dependence of the differential multiplicity of neutral pions is shown. This quantity is defined as the number $\langle N^\pi \rangle$ of pions produced in the deep-inelastic reaction normalised to the total number $\langle N_{\text{DIS}} \rangle$ of inclusive events [126]

$$\frac{1}{N_{\text{DIS}}(Q^2)} \frac{dN^\pi(z,Q^2)}{dz} = \frac{\sum_q e_q^2 \int_{0}^{1} q(x,Q^2) D^\pi_q(z,Q^2) dx}{\sum_q e_q^2 \int_{0}^{1} q(x,Q^2) dx}.$$  

(67)

The data are compared to results from SLAC [127] and from EMC [128], taken at 10-times higher muon beam energy. The HERMES results are systematically larger than those from EMC. This difference can, however, be totally explained by the different $Q^2$ ranges of the two experiments: $\langle Q^2 \rangle = 2.5 \text{ GeV}^2$ for HERMES and $\langle Q^2 \rangle = 25 \text{ GeV}^2$ for EMC. The HERMES data have been evolved to the mean $Q^2$ of the EMC data using a NLO-QCD model for the evolution of fragmentation functions [121]. The result is shown in the right panel of Fig. 23. There is very good agreement within the uncertainties. As a consequence one can conclude that either the data at 10-times higher beam energy suffer from the same problems mentioned above in section 4.4.1 or that such a semi-inclusive analysis can be safely performed at HERMES energies as well. Also the multiplicities for neutral pions and the average multiplicities for charged pions agree within errors over a large range of $z$. This is shown in Fig. 24, where their ratio is seen to be consistent with unity below $z \sim 0.70$. At larger $z$ the multiplicity is larger for charged pions than for neutral pions. This might be related to hard exclusive electro-production [129] or to higher twist processes. The inclusion of events from the high $z$ region in the analysis of spin-dependent quark distributions may therefore need special consideration.

![Figure 23: Multiplicity of neutral pions from a hydrogen target as a function of $z$ at the measured values of $Q^2$ (left panel) and after evolution of the HERMES data to the mean $Q^2$ value of the EMC data (right panel).](image)

### 4.4.3 Flavour asymmetry of the light quark sea

The flavour asymmetry of the light quark sea, i.e. $\bar{d} - \bar{u}$, has been determined at HERMES by combining the multiplicities for negative and positive pions from unpolarised hydrogen and deuterium
targets [130]. In Fig. 25 (left) the difference $\bar{d} - \bar{u}$ is shown in comparison to the data from the Drell-Yan experiment E866 [131, 132]. There is good agreement within the accuracy of the HERMES data. It has also been demonstrated that for fixed values of $x$ the ratio $(\bar{d} - \bar{u})/(u - d)$ does not depend on $z$ within errors, see Fig. 25 (right). The $z$-independence of the ratio $(\bar{d} - \bar{u})/(u - d)$ provides further evidence in support of the factorisation ansatz expressed by Eq. (64). Apparently, target fragmentation plays only a minor role in the kinematic configuration of the HERMES experiment. For a recent review of this exciting subject see Ref. [133].
4.4.4 Quark polarisations and purities

The quark polarisations $\Delta q(x, Q^2)/q(x, Q^2)$ are isolated in LO QCD by introducing purities $P_q^h$ [134], which are unpolarised quantities and defined by:

$$P_q^h (x_i) = \frac{\epsilon_{q}^2 (x_i, Q^2) \mathcal{D}_q^h (x_i, z_i, Q_i^2)}{\sum_{q'} \epsilon_{q'}^2 (x, Q^2) \mathcal{D}_{q'} (x_i, z_i, Q_i^2)} ,$$

where $\mathcal{D}_q^h (x, z_i, Q_i^2)$ are effective fragmentation functions in a bin $i$ with mean values of the kinematical quantities $x_i$, $z_i$, and $Q_i^2$ derived by averaging over all events in the bin, which take into account also the acceptance function of the spectrometer. Note that the purities only depend on one single parameter $x_i$ in this definition, as the mean values $Q_i^2$ and $z_i$ are fixed within one $x$-bin $i$.

These purities may be interpreted as the probability that, in a given bin of $x$, the photon struck a quark of type $q$ in the nucleon when a hadron of type $h$ is detected in the experiment. Obviously, the purities for each hadron species add up to one in each bin.

With this definition of the purities Eq. (65) can be rewritten as

$$A^h (x_i) = C_i \sum_q P_q^h (x_i) \frac{\Delta q (x_i, Q_i^2)}{q (x_i, Q_i^2)} .$$

The purities provide a simple way to separate the spin-dependent quark distributions from other quantities, which are related to unpolarised quark distributions and fragmentation functions. Their input parameters are known with good precision from a large number of DIS experiments on unpolarised targets, augmented by data on fragmentation functions from $e^+e^-$ annihilation experiments.

For the determination of the various quark polarisations one combines inclusive asymmetries and semi-inclusive asymmetries for positive and negative hadrons and for identified hadrons of different species from various targets and obtains a system of linear equations which can be written for each bin $i$ in a matrix form as:

$$\bar{A} (x_i) = C_i \cdot \mathbf{P} (x_i) \bar{Q} (x_i) .$$

Here the vector $\bar{A}$ contains as elements the measured inclusive and semi-inclusive hadron asymmetries for the various targets, the vector $\bar{Q}$ contains the quark and anti-quark polarisations for each flavour and $\mathbf{P}$ is the purity matrix with elements $P_q^h$. This system of asymmetries is typically over-determined, as the number of measured asymmetries is larger than the number of various quark polarisations to be extracted.

The polarisations for each bin $x_i$ are then obtained by inverting Eq. (70) and minimising

$$\chi^2 = (\bar{A} - \mathbf{P} \cdot \bar{Q})^T \nu^{-1} (\bar{A} - \mathbf{P} \cdot \bar{Q}) ,$$

where $\nu^{-1}$ is the covariance matrix of the asymmetry vector $\bar{A}(x)$.

4.4.5 Spin-dependent quark distributions from hydrogen, $^3$He and deuterium

In a first analysis HERMES has used the data from the $^3$He and the hydrogen targets taken in 1995 and 1996/97 [135]. The kinematic requirements for the inclusive data used in this analysis were $Q^2 > 1$ GeV$^2$, $y < 0.85$, $W^2 > 4$ GeV$^2$. Additional requirements were imposed on the coincident hadrons to determine the semi-inclusive asymmetries. As the RICH information was not yet available for these periods of data taking, the semi-inclusive asymmetries for the different hadron species were combined into those for positive hadrons $h^+$ and negative hadrons $h^-$. To select hadrons related to the struck quark a minimum $z$ of 0.2 and $x_F \approx 2p_L/W$ of 0.1 were required, where $p_L$ is the longitudinal momentum of the hadron with respect to the virtual photon in the photon-nucleon centre-of-mass system. Also a minimum of $W^2$ of 10 GeV$^2$ was used in order to improve the separation of the current and target fragmentation regions, and thus increase the sensitivity to the struck quark. After applying all cuts about 2 million inclusive DIS events remained for both targets, about 13% of those accompanied by a
positive hadron and 7% by a negative hadron. It should be noted here that for the second analysis, which also includes the large statistics sample from the deuterium target taken in 1998-2000, the $W^2$ cut for the inclusive sample has been increased to 10 GeV$^2$, like in the semi-inclusive case, and the $z$-range was reduced from $0.2 < z < 1.0$ to $0.2 < z < 0.8$ to discard any hadrons that originate from exclusive processes. All asymmetries were corrected for detector smearing and charge symmetric background. QED radiative corrections were only applied to the inclusive asymmetries, those on the semi-inclusive asymmetries being negligible. The dominant sources of systematic uncertainties are due to the target and beam polarisation measurements.

![Figure 26: The inclusive (left) and semi-inclusive asymmetries for positively (middle) and negatively (right) charged hadrons for a proton (top) and a $^3$He (bottom) target.](image)

Fig. 26 shows the extracted inclusive and semi-inclusive asymmetries for positively and negatively charged hadrons on both targets. The measured spin asymmetries $A_\Sigma^\pm(x, Q^2, z)$ were integrated in each $x$ bin over the corresponding $Q^2$ range and over the $z$ range from 0.2 to 1.0 to yield $A_\Sigma^\pm(x)$. Also shown in this figure are the inclusive asymmetries measured at a similar energy at SLAC [136, 73, 77, 82], and semi-inclusive hadron asymmetries on hydrogen measured by SMC [137]. The data are in agreement within the quoted uncertainties. The agreement with the SMC data, which are taken at a typical average $Q^2$ value 6-12 times higher compared to the HERMES data, illustrates that the semi-inclusive asymmetries are independent of $Q^2$ within the present accuracy. Although the statistics of the $^3$He sample is similar to that of the hydrogen sample, the statistical error on the asymmetries is much larger due to the smaller dilution factor $f \approx 1/3$, and the smaller target polarisation $P_T \approx 0.46$.

The calculated purities include the effects of the acceptance of the experiment. They have been determined with a Monte Carlo Simulation using the LUND string fragmentation model [138], a GEANT based model of the detector, the CTEQ Low-$Q^2$ parameterisation [139] for the unpolarised parton distributions, and values of $R$ from Ref. [67]. The LUND fragmentation parameters were tuned to fit the measured hadron multiplicities. The generation of these purities is explained in detail in several HERMES PhD Theses (see for example [140, 141, 142]). As an example, the purities for a proton target [142] are shown in Fig. 27. It can be seen that even in the case of negative hadrons, which in general do not contain an up quark as constituent, the purities are dominated by the unfavoured fragmentation of
up quarks, since $e_u^2 = 4 \cdot e_d^2$ and $u(x) \approx 2 \cdot d(x)$. It is especially worth noting here that this is also true \[141\] in the case of the $K^- = (\bar{u}s)$, which is a pure sea-quark object.

In principle the six measurements displayed in Fig. 26 should allow to determine six quark and anti-quark polarisations. To improve the statistical significance only three polarisations have been fitted, imposing constraints on the sea polarisation. For most of the analysis it was assumed that the sea polarisation $\Delta q_s(x)/q_s(x)$ is independent of flavour:

\[
\begin{align*}
\Delta u_s(x) &= \Delta d_s(x) = \Delta s_s(x) = \Delta \bar{u}(x) = \Delta \bar{d}(x) = \Delta \bar{s}(x) = \Delta q_s(x) = q_s(x) 
\end{align*}
\]

(72)

In order to enable a comparison with the data from SMC their assumption on the polarised sea distributions was also adopted:

\[
\Delta u_s(x) = \Delta d_s(x) = \Delta s_s(x) = \Delta \bar{u}(x) = \Delta \bar{d}(x) = \Delta \bar{s}(x). 
\]

(73)

The fitted quark polarisations differ by typically less than 0.01 if either one of the two assumptions is used.

The polarisation vector $\vec{Q}$ has been chosen as

\[
\vec{Q} = \left( \frac{\Delta u(x) + \Delta \bar{u}(x)}{u(x) + \bar{u}(x)}, \frac{\Delta d(x) + \Delta \bar{d}(x)}{d(x) + \bar{d}(x)}, \frac{\Delta s(x) + \Delta \bar{s}(x)}{s(x) + \bar{s}(x)} \right), 
\]

(74)

where due to the assumption expressed by Eq. (72) the polarisation of the strange quarks and the total sea are equal

\[
\frac{\Delta s(x) + \Delta \bar{s}(x)}{s(x) + \bar{s}(x)} = \frac{\Delta q_s(x)}{q_s(x)}. 
\]

(75)

For $x > 0.3$ the sea polarisation is set to zero, and the corresponding effect on the results for the non-sea polarisations is included in their systematic uncertainties.

Fig. 28 (left) shows the results. The up quark polarisations are positive and the down quark polarisations are negative over the measured range of $x$. Their absolute values are largest at large $x$ and remain different from zero at low values of $x$, where also the sea quarks contribute. The sea polarisation is compatible with zero over the measured range of $x$.

The spin-dependent quark distributions $\Delta q(x)$ were determined by taking the products of the three polarisations and the corresponding spin-independent quark distributions from Ref. [139] at $Q^2 = 2.5$ GeV$^2$. The results for the up and down distributions are shown in the right panel of Fig. 28 and
Figure 28. The extracted quark polarisations as a function of $x$ (left panel). In the right panel $x(\Delta u(x) + \Delta \bar{u}(x))$ and $x(\Delta d(x) + \Delta \bar{d}(x))$ are shown as a function of $x$ at $Q^2 = 2.5$ GeV$^2$. They are compared to the parameterisations of Refs. [91, 74, 96]. The parameterisations from Refs. [91, 96] are corrected by a factor $(1 + R)$ to allow for a direct comparison.

compared with different sets of parameterisations of world data in LO QCD [143, 93, 115]. At the lowest $x$ values the absolute magnitude of both distributions is about equal, while $\Delta d(x)/\Delta u(x) \approx -0.4$ at $x$ values around 0.2, which is consistent with the negative value of $q^u_s(x)$. While $x(\Delta u(x) + \Delta \bar{u}(x))$ is already very well determined from the data, the errors for $x(\Delta d(x) + \Delta \bar{d}(x))$ are still rather large. This is mainly due to the relatively poor statistical accuracy of the $^3$He asymmetries.

The situation improves substantially when the high statistics deuterium data from 1998-2000 are included in the analysis. From these data semi-inclusive asymmetries could not only be determined for positively and negatively charged hadrons but also for identified positive and negative pions and kaons. The analysis is still in progress. Fig. 29 (left panel) shows the preliminary semi-inclusive asymmetries for positive and negative hadrons from the deuteron target compared with those from SMC [137]. The enormous statistical improvement compared to the $^3$He data and also the SMC data is evident. The much bigger data sample allowed to drop the constraints on the sea polarisation (72) and to perform a five parameter fit, using the following vector $\vec{Q}$: $\vec{Q} = (\Delta u_s(x), \Delta d_s(x), \Delta d_s(x), \Delta d_s(x), \Delta s_s(x), \Delta s_s(x))$. The preliminary results for the spin-dependent quark distributions are shown in the right panel of Fig. 29. The statistical accuracy of the polarised distributions has improved significantly. However, because the radiative and smearing effects are unfolded, rather than corrected for, the improvement of the accuracy is less than could have been expected on the basis of the statistical improvement alone. All measured spin-dependent sea-quark distributions are compatible with zero, and also compatible with the recent parameterisations from QCD fits to the inclusive data [115, 144]. In these calculations the assumption of a $SU(3)_f$ symmetric polarised sea results in a small and slightly negative polarised sea
quark distribution, on the other hand the measured strange sea polarisation has a hint of being slightly positive at low $x$. More precise data are needed to assess whether a true discrepancy occurs in this sector.

4.5  The spin-dependent gluon distribution $\Delta g(x)$

4.5.1 Methods to determine the spin-dependent gluon distribution

As discussed above, the apparent deficit in the contribution of quark spins to the nucleon spin can be explained by assuming a substantial gluon polarisation. The question is how this assumption can be investigated experimentally. The determination of the gluon distribution is a rather difficult task since the virtual photon does not couple directly to the gluon.

The unpolarised gluon distribution $g(x, Q^2)$ is rather well known from global QCD fits [145, 146, 103] of the unpolarised analogue of the DGLAP equations (51,52) to the $x$ and $Q^2$ dependence of the unpolarised deep-inelastic scattering cross sections or structure functions respectively, which are experimentally known with high precision from $e$-p collider and fixed target experiments [87, 88, 89, 90] over a large range in $x$ and $Q^2$. These global fits include also experimental information from other processes where the cross section is proportional to the gluon density like lepton-pair, di-jet or direct photon production.

In inclusive polarised deep-inelastic scattering the situation is very different. As shown in Sect. 4.2, $\Delta g(x, Q^2)$ is essentially unconstrained by the present data. Even if additional more precise fixed target data become available, as expected mainly from COMPASS, it is rather unlikely that inclusive data will sufficiently constrain the spin-dependent gluon distribution. The reason is the very weak $Q^2$ dependence of the inclusive spin asymmetries. The situation would change, however, with the realisation of a polarised HERA $\bar{e}p$ collider, since such a facility would cover a much larger range in $x$ and $Q^2$, and thus enable to probe the spin-dependent gluon distribution through the $Q^2$ evolution of the $g_1$ structure function.

Therefore, at the moment the most promising approach for studying the spin-dependent gluon distribution seems to be direct measurements via processes where the cross section is proportional directly
to the gluon density or to the square of the gluon density. In lepton-nucleon scattering such a process is photon-gluon fusion (PGF) [147, 148], depicted in Fig. 30, where the virtual photon and a gluon from the target interact by the exchange of a virtual quark (anti-quark). The result is a quark and anti-quark pair which is produced back-to-back in the photon-gluon centre-of-mass system. This process can be isolated either by studying reactions in which a charm quark and an anti-charm quark are produced and the necessary hard scale is given by (2 times) the mass of the charm quark. Alternatively the PGF process can be identified by measuring a quark and an anti-quark jet at large transverse momenta $p_t$ relative to the direction of the virtual photon [149, 150]. Here $p_t$ has to be at least of the order of 1 GeV to ensure the hard scale.

Charm production via photon-gluon fusion can be tagged by detecting $J/\Psi$ production (hidden charm), or by the decay of a $D$ or $D^*$ meson which originated from the fragmentation of one of the charm quarks (open charm). Indeed data for hidden charm production from the fixed target muon experiments EMC and NMC [151, 152, 153] and the collider experiments H1 and ZEUS [154, 155, 156] and also for open charm production [157, 158] have been used to determine the unpolarised gluon distribution. It has also been extracted at high $Q^2$ and low values of $x$ by the collider experiments ZEUS and H1 [159] from di-jet events with large transverse momentum, which originate dominantly from photon-gluon fusion and the QCD Compton process.

The HERMES spectrometer has been upgraded by several detectors (an instrumented iron wall behind the calorimeter, the RICH detector, the Forward Quadrupole Spectrometer, and hodoscopes before and after the HERMES magnet outside the standard acceptance) to enhance the acceptance for particles originating from the decays of charmed particles and the capabilities to identify them. The experimentally most easily accessible channel would be the decay $D \rightarrow \pi K$. This decay channel suffers, however, from a large background from non-resonant $\pi K$ production, which can be suppressed by demanding a soft pion from the decay $D^* \rightarrow D\pi_{\text{soft}}$ among the final state hadrons. Unfortunately the HERMES centre-of-mass energy is just a few GeV above the charm production threshold where the cross section is very small. Therefore, the statistics for events from charm photoproduction cross section, and the corresponding double spin asymmetries can be found in [160, 161]. These analyses showed that the presently available statistics is unsufficient to measure $\Delta g$.

The perspectives are much better for the COMPASS [162] experiment, which in spring 2002 has started data taking with a muon beam energy of 160 GeV (where the charm cross section is much higher). It is expected that ultimately the error on the asymmetry will be as low as 0.04. Very high precision could be obtained by the proposed SLAC experiment E161 [163], which intends to determine the spin-dependent gluon distribution via open charm photo-production with a polarised real photon beam, generated by coherent bremsstrahlung from a diamond crystal. The authors of this proposal
claim that for each of the three peak energies of 35, 40 and 45 GeV a statistical accuracy of 0.006 in the asymmetry for charm production can be achieved.

At a polarised HERA $\bar{e}p$-facility di-jet production could be a useful process to study the polarised gluon density [164, 165, 166, 167] down to $x$ values of $10^{-3}$. At RHIC $\Delta g(x)$ will be measured via direct photon and di-jet production in polarised proton-proton scattering [168] with a precision similar, or even slightly better than expected for the COMPASS experiment but an entirely different systematic uncertainty.

4.5.2 HERMES results for the spin asymmetry in the photo-production of pairs of high-$p_t$ hadrons

At the relatively low centre-of-mass energy of polarised fixed target experiments only few hadrons are produced and observed. It has been proposed by Bravar, Kotzinian and von Harrach [169] to use in this case pairs of correlated hadrons at large transverse momenta instead of di-jets to tag PGF. Several calculations for large-$p_t$ hadron-pair production for HERMES and COMPASS kinematics [170, 171] confirmed the feasibility of this approach. Because of the helicity selection rule the photon and the gluon must have opposite helicities in the photon-gluon fusion process to produce a quark and anti-quark and hence the analysing power for the PGF subprocess is $-1$ (for massless quarks). Hence, if a negative double spin asymmetry for pairs of correlated hadrons with high transverse momenta is measured on a polarised proton target one can safely conclude that the gluon polarisation is positive since all other possible subprocesses contributing to the asymmetry have a positive analysing power. To extract the magnitude of the polarised gluon distribution from such an asymmetry one needs, however, a good knowledge of the relative contribution of all subprocesses.

Following the ideas and the kinematic requirements suggested in [169], HERMES has studied the asymmetry for the production of pairs of hadrons at high transverse momenta [172]. Hadrons of opposite charge have been selected as the product of two favoured fragmentation functions then enters the asymmetry. Events were selected that contained at least one positively charged hadron $h^+$ and at least one negatively charged hadron $h^-$. The hadrons were required to have a momentum $> 4.5$ GeV and a transverse momentum $p_t > 0.5$ GeV relative to the positron beam direction. A minimum value of the two-hadron invariant mass (assuming both hadrons to be pions) was imposed. The observation of the scattered positron was not required in the trigger. Therefore, the data sample is dominated by events from very low $Q^2$ and the average kinematics of the data sample had to be determined from a Monte Carlo simulation. The selected data sample from the proton target, collected in 1996-1997 contained about 600 events. Fig. 31 (left panel) shows the measured asymmetry $A_1$ for the highest transverse momenta accessible at HERMES. The distribution is limited to $p_t < 2$ GeV as the beam energy of 27.5 GeV and the forward angular acceptance of < 220 mrad do not give access to higher $p_t$ values. The fastest hadron $h_1$ was required to have a $p_t > 1.5$ GeV. $A_1$ is then plotted as a function of the transverse momentum, $p_t^{h_2}$, of the hadron with opposite charge and second highest transverse momentum. For $p_t^{h_2} > 1$ GeV the double spin asymmetry is found to have a negative sign: $A_1 = -0.28 \pm 0.12\text{ (stat.)} \pm 0.02\text{ (sys.)}$.

The observed negative asymmetry is in contrast to the positive asymmetries typically measured in DIS from protons and, as explained above, is an indication for a positive gluon polarisation.

To relate the measured negative asymmetry to the gluon polarisation one has to understand precisely the relative contribution of various subprocesses to the hadron pair production cross section at HERMES kinematics. Therefore a Monte Carlo simulation in leading order has been performed, using version 5.724 of the PYTHIA generator [173] and version 7.410 of JETSET. Several subprocesses can contribute to the two-hadron photoproduction cross section: photon-gluon fusion (PGF), the QCD Compton effect (QCDC), where the struck quark radiates a hard gluon, interaction via the hadronic structure of the photon described by the vector meson dominance model (VMD), non-resonant hadronic ‘anomalous’ photon structure and finally usual hadron production in lowest order DIS. The contribution from the latter process is suppressed by the requirement of high $p_t$, and also a negligible contribution from anomalous photon structure is predicted at HERMES energies. Contributions from VMD were assumed to have a negligible spin asymmetry and were treated as a dilution to the other asymmetries.
Figure 31: $A_{\parallel}$ for pairs of high-$p_t$ hadrons measured at HERMES from a polarised proton target compared with Monte Carlo predictions for different assumptions about the magnitude of the gluon polarisation (left), Monte Carlo prediction for the fraction of events coming from each relevant process type: PGF (solid line), QCDC (dashed) and VDM (dotted) (middle), and the HERMES result for the gluon polarisation $\Delta g/g(x_G)$ from the asymmetry for pairs of high-$p_t$ hadrons. The latter result is compared to LO-QCD fits to a subset of the world’s data on $g_1$ (right).

Under these assumptions only $A_{\text{PGF}}$ and $A_{\text{QCDC}}$ contribute significantly to the measured asymmetry:

$$A_{\parallel} \approx (A_{\text{PGF}} f_{\text{PGF}} + A_{\text{QCDC}} f_{\text{QCDC}}) D \approx (\hat{A}_{\text{PGF}} \frac{\Delta g}{g} f_{\text{PGF}} + \hat{A}_{\text{QCDC}} \frac{\Delta q}{q} f_{\text{QCDC}}).$$

(76)

The subprocess asymmetries $\hat{A}_{\text{PGF}}$ and $\hat{A}_{\text{QCDC}}$ are calculable in leading order QCD. For real photons and massless quarks $\hat{A}_{\text{PGF}} = -1$ and $\langle \hat{A}_{\text{QCDC}} \rangle \approx +0.5$ (averaged over the selected kinematics) independently of quark flavour. The effective quark polarisation $A_q/q$ is calculated as the weighted mean of the known up and down quark polarisations $A_u/u$ and $A_d/d$.

The unpolarised fractions $f_i$ of the events from the three relevant subprocesses $i$ ($f_{\text{PGF}} + f_{\text{QCDC}} + f_{\text{VMD}} = 1$), as determined from the Monte Carlo simulation, are plotted in Fig. 31 (middle panel) vs. $p_t^H$. In the region of phase space where a negative asymmetry is observed the simulated cross section is dominated by PGF.

With these assumptions and model parameters and after averaging over a relatively small region of phase space the gluon polarisation $\Delta g/g$ at the average kinematics of the experiment has been determined. First the central kinematical parameters were extracted from the Monte Carlo simulation, yielding $\langle y \rangle = 0.86, \langle D \rangle = 0.93, \langle D_{\text{QCDC}} \Delta q \rangle = 0.15, \langle Q^2 \rangle = 0.06 \text{ GeV}^2$ and $\langle p_t^H \rangle = 2.1 \text{ GeV}$

Moreover the measured double spin asymmetry has been extracted by solving Eq. (76):

$$\langle \Delta g/g \rangle = 0.41 \pm 0.18 \text{ (stat.)} \pm 0.03 \text{ (exp.sys).}$$

As shown in Fig. 31 (right panel), where the extracted $\Delta g/g$ is compared with several model curves, the distribution $\Delta g(x_y)$ is probed in the region $0.05 < x_y < 0.28$, with the average value $\langle x_y \rangle = 0.17$. The sensitivity of the asymmetry to the magnitude of the spin-dependent gluon distribution is shown by the different curves in the left panel of Fig. 31.

This result allows the first direct glimpse on the spin-dependent gluon distribution. As pointed out above the extracted value depends a lot on several assumptions made about the relevant subprocesses and on the applicability and reliability of the PYTHIA generator at the relatively low energy of the HERMES experiment. These uncertainties might be rather large, but hard to estimate. At the kinematics of the experiment, no spin dependent analysis of higher order QCD processes is available, for example. To alter, however, the principle conclusion that the gluon polarisation is positive in the covered range of $x_y$, a significant contribution from a neglected process with large negative spin asymmetry would be needed, which is rather unlikely.

In the years 1998-2000 about 9 million DIS events were collected from the polarised deuterium target as compared with 2.4 million from hydrogen. Performing the same analysis as in the hydrogen case, the asymmetry $A_{\parallel}$ for pairs of high-$p_t$ hadrons was also determined. The availability of the RICH allowed to do a separate analysis for identified pion and kaon pairs. As strange quarks are suppressed relative...
to up and down quarks in the nucleon target, in the fragmentation and in the hadronic fluctuations of
the photon, correlated high-\textit{p}_T kaon pairs should be especially sensitive to PGF. The selection of kaons,
however, reduces the event yield such that the present statistics does not yet improve the significance
of the measurement. A new extraction of \((\Delta g/g)\) based on all HERMES data is currently underway
using improved simulations.

The double spin asymmetry in the production of pairs of high-\textit{p}_T hadrons and especially kaons will,
in addition to the asymmetry in \(D^0\) and \(D^*\) production, be used by the COMPASS experiment to
determine \(\Delta g/g\) with much better statistical accuracy.
5 Hard exclusive reactions and generalised parton distributions

5.1 Generalised parton distributions

Recently it has been shown [174] that the total angular momentum of quarks

\[ J^q = \frac{1}{2} \Delta \Sigma + L_z^q \]  

(77)

can be related to the second moment of certain Generalised Parton Distributions (GPDs) which can be accessed in electroproduction processes by studying hard exclusive reactions. But it is also believed that these distributions provide access to certain non-perturbative aspects of the QCD structure of the nucleon. This is the main reason for the renewed and growing interest in these reactions.

In electroproduction processes hard exclusive reactions of two types are discussed:

\[ \gamma^*(q) + T(p) \rightarrow \gamma(q') + T'(p') \]

\[ \gamma^*(q) + T(p) \rightarrow M(q') + T'(p') \]  

(78)

in which a virtual photon \( \gamma^* \) with high energy and large virtuality \( -q^2 = Q^2 > 0 \) interacts with a hadronic target \( T \) and produces a real photon \( \gamma \) or meson \( M \), and a low-mass hadronic state \( T' \) (see Fig. 32).

![Handbag diagrams for DVCS (left) and for exclusive meson production (right).](image)

Figure 32: Handbag diagrams for DVCS (left) and for exclusive meson production (right).

It was realised that these hard reactions can be interpreted in a QCD framework when it was shown [175, 176, 177, 178] that a clear separation into a perturbative and a non-perturbative part of the interactions could be realized. Moreover, the amplitudes of the processes (78) can be factorised into three terms: i) a hard scattering coefficient which describes the short distance perturbative stage of the interaction and can be calculated in NLO QCD [177, 179, 180, 181], ii) a meson distribution amplitude, which describes how the meson is formed from the two partons and thus contains non-perturbative meson physics, and iii) a Generalised Parton Distribution (GPD) [182, 183, 184, 174, 185, 175], which contains non-perturbative nucleon structure information. It is these GPDs which are of importance in the context of this review. Due to the factorisation theorem the same universal distributions enter the description of various hard reactions. They provide a unified formalism for the description of inclusive deep-inelastic scattering, exclusive meson production, deeply-virtual Compton scattering (DVCS) and electromagnetic form factors. Furthermore, as mentioned above, the second moment of certain GPDs can be related to the total angular momentum \( J^q \) carried by the quarks in the nucleon. Using the information on \( \Delta \Sigma \) available from inclusive and semi-inclusive polarised deep-inelastic scattering, this may be used to derive (gauge dependent) information on the quark orbital momentum contribution \( L_z^q \) to the spin of the nucleon. At present this seems to be the best way to access this quantity experimentally. To my knowledge there exists only one other suggestion how to measure \( L_z^q \). This simple qualitative model [186] relates 'rotating constituents' in a tranversely polarised nucleon to an azimuthal asymmetry of the produced hadrons relative to the target polarisation plane.

For details the reader is refered to the recent excellent review in Ref. [187] in which the physics content of the GPDs is highlighted and it is discussed how they can be accessed from a wide variety of
hard electroproduction processes. Here only the main properties of the GPDs are summarised, which are important for the results presented below.

At leading twist-2 level, there are four (quark chirality conserving) different types of quark GPDs: the unpolarised distributions \( H^q \) and \( E^q \) and the polarised distributions \( \tilde{H}^q \) and \( \tilde{E}^q \). The GPDs \( H^q \) and \( \tilde{H}^q \) conserve nucleon helicity, while \( E^q \) and \( \tilde{E}^q \) are associated with a helicity-flip of the nucleon. The GPDs depend on three variables \( \hat{x}, \xi \) and \( t \). The variable \( \hat{x} \) is the light cone momentum fraction defined by \( k^+=\hat{x}P^+ \), where \( k \) is the quark loop momentum in Fig. 32, and \( P \) is the average nucleon momentum \( \frac{p+p^'}{2} \). (Note that in the literature commonly \( x \) is used for this light cone momentum fraction. In this review it is consistently replaced by \( \hat{x} \) to avoid confusion with Bjorken \( x \)). The four-momentum difference \( \Delta = p'-p \) is the overall momentum transfer and \( t = \Delta^2 \) is the total squared momentum transfer. The third variable entering the GPDs is the skewedness variable \( \xi \) which is defined by \( \Delta^+ = -2\xi P^+ \). In the Bjorken limit \( \xi = x/(2-x) \). The GPDs describe quark-quark correlations where a quark with momentum fraction \( \hat{x}-\xi \) is taken out of the initial nucleon (having momentum \( p \)) and put back into the final nucleon (having momentum \( p^' \)) with momentum fraction \( \hat{x}+\xi \).

In the forward limit \( t \to 0, \xi \to 0 \)
\[
H^q(\hat{x},0,0) = q(\hat{x}), \quad \tilde{H}^q(\hat{x},0,0) = \Delta q(\hat{x}),
\]  
the ordinary quark and quark helicity distributions \( q(x) \) and \( \Delta q(x) \) are obtained. The variable \( \hat{x} \) is defined in the range \((-1,+1)\), where negative values of \( \hat{x} \) correspond to anti-quark distributions in the following manner:
\[
q(-\hat{x}) = -\bar{q}(x), \quad \Delta q(-\hat{x}) = \Delta \bar{q}(x).
\]  
The functions \( E \) and \( \tilde{E} \) are not accessible in deep-inelastic scattering, but hard exclusive production processes are dependent on \( E \) and \( \tilde{E} \).

The first moments of the GPDs are related for any \( \xi \) to the elastic nucleon form factors through model independent sum rules [174]:
\[
\sum_q e_q \int_{-1}^1 d\hat{x} H^q(\hat{x},\xi,t) = F_1(t), \quad \sum_q e_q \int_{-1}^1 d\hat{x} E^q(\hat{x},\xi,t) = F_2(t),
\]
\[
\sum_q e_q \int_{-1}^1 d\hat{x} \tilde{H}^q(\hat{x},\xi,t) = g_A(t), \quad \sum_q e_q \int_{-1}^1 d\hat{x} \tilde{E}^q(\hat{x},\xi,t) = h_A(t),
\]
where \( F_1(t) \) and \( F_2(t) \) are the elastic Dirac and Pauli nucleon form factors, and \( g_A(t) \) and \( h_A(t) \) are the axial-vector and pseudo-scalar form factors.

Finally there is the relation between the second moment of a combination of GPDs with a given flavour \( q' \) on the one hand, and the total angular momentum \( J^{q'} \), i.e. the sum of intrinsic and orbital angular momenta carried by quarks of the flavour \( q' \) in the proton [174], on the other hand:
\[
\lim_{t \to 0} \frac{1}{2} \int_{-1}^{+1} d\hat{x} \hat{x} [H^{q'}(\hat{x},\xi,t) + E^{q'}(\hat{x},\xi,t)] = J^{q'}.
\]
A determination of the contributions \( J^{q'} \) from all quark flavours gives the total quark angular momentum \( J^q \).

So the orbital angular momentum part \( L^z \) can in principle be accessed, if one has determined the GPDs \( H \) and \( E \) in the limit \( t \to 0 \). Due to the multi-dimensional dependence of the GPDs (also \( Q^2 \) enters the dependence), the GPDs cannot be easily unfolded from actual data. The study of many different processes will be required for the determinations of the GPDs. Exclusive vector meson production, for instance, is only sensitive to the distributions \( E \) and \( H \), pseudo-scalar meson production to the distributions \( E \) and \( \tilde{H} \). Deeply virtual Compton scattering (DVCS) on the other hand is sensitive to all four distributions. Based upon today’s knowledge of QCD, and similar to the PDFs, GPDs cannot be calculated from first principles and up to now no measurements of GPDs exist.

It appears worth noting that the measurement of hard exclusive processes is the main scientific objective of a long-term-future project for a high-luminosity successor of the HERMES experiment [188].
5.2 Deeply Virtual Compton Scattering

5.2.1 Kinematics and asymmetries

A particularly promising way of studying the nucleon structure is through deeply virtual Compton scattering (DVCS), where a virtual photon couples to the proton, and a real photon is emitted ($\gamma^* + p \rightarrow \gamma + p$), as illustrated in Fig. 33 (left panel). The proton is left intact. The DVCS process combines the advantages of deep inelastic scattering (DIS) and Compton scattering. The initial photon is off-shell, thus opening an additional dimensional parameter with respect to real Compton scattering, because the virtuality can be 'tuned' (in the given kinematic limits). This reaction is unique, as the real photon carries direct information about the partonic structure of the nucleon, there is no non-perturbative meson distribution amplitude, and the reaction can therefore be described with GPDs only.

![Diagram of DVCS and Bethe-Heitler processes](image)

Figure 33: Diagrams for the DVCS (left-a) and the Bethe-Heitler processes (right-b).

In electron (or positron) scattering experiments there exists another process that leads to the same final state as DVCS, namely the Bethe-Heitler process (BH), in which the incoming or outgoing electron radiates a Bremsstrahlung photon in the Coulomb field of the proton, see Fig. 33 (right panel). Hence, in the process $e p \rightarrow e p \gamma$ the amplitudes of DVCS and BH are interfering. The cross section for the exclusive leptoproduction of real photons reads

$$\frac{d^4\sigma}{d\phi dtdQ^2 dx} = \frac{xy^2}{32(2\pi)^4 Q^4} \frac{|A_{BH} + A_{DVCS}|^2}{(1 + \gamma^2)^{1/2}},$$

(83)

where $A_{BH}$ and $A_{DVCS}$ represent the amplitudes of the BH and DVCS processes.

The relative contributions of BH and DVCS to the total BH/DVCS amplitude depend strongly on the lepton energy. At HERMES energies, the BH contribution is mostly dominant, except at the largest values of $Q^2$ and the Bjorken variable $x$. The dominance of the BH contribution is used as an opportunity since the interference of the two amplitudes offers a way to access both the real and the imaginary part of the DVCS amplitude (see below). At the high energies available at H1 and ZEUS, where the DVCS amplitude dominates over the BH amplitude, the square of the DVCS amplitude is accessed.

The amplitudes describing the DVCS and BH processes, and their interference can be expanded into a $1/Q$ power series. All terms in the expansion depend (each in a different way) on the azimuthal angle $\phi$ between the scattering plane and the production plane (see Fig. 34), and can be expressed in moments $\langle \cos n\phi \rangle$ and $\langle \sin n\phi \rangle$. These cosine and sine moments are sensitive to the real and imaginary parts of the DVCS helicity amplitudes, respectively. In order to project out the various moments, measurements have to be done with different beam helicity and charge. Assuming the validity of the handbag diagram (Fig. 33, left panel), each moment has its own characteristic fall-off with $(1/Q)^n$ at fixed $x$ and $t$. The first cosine moment is then dominant with a $1/Q$ fall-off, while higher moments fall off with powers of 2 or higher [189].

A theoretical picture of the DVCS process has been developed in a number of papers [174, 190, 189, 191]. Explicit expressions for the amplitudes of the DVCS, Bethe-Heitler and interference terms including the first sub-leading correction in $1/Q$ for different kinds of polarised and unpolarised initial particles have been calculated in Refs. [192, 193, 194, 195].
Figure 34: Schematic visualisation of the azimuthal angle $\phi$ between the scattering and the production plane. The definition of $\phi$ is equally valid for DVCS and meson production.

To access, as a first step, the nucleon helicity conserving GPD $H(\hat{x}, \xi, t)$ an unpolarised proton target is sufficient. In this case two different types of measurements are possible:

i) the lepton beam-helicity asymmetry, which is (in the leading twist-2 approximation) proportional to the difference of the cross sections for scattering of longitudinally (L) polarised positron (or electron) beams with opposite helicities off an unpolarised (U) target. The cross section difference is dependent on $\sin \phi$.

$$d\Delta \sigma_{LU} \equiv d\sigma(e^+p) - d\sigma(e^-p) \sim \mp \sin \phi \times \text{Im} \mathcal{A},$$  \hspace{1cm} (84)

ii) the lepton charge asymmetry, which is proportional to the difference of the cross sections for scattering of unpolarised leptons of either charge off an unpolarised target. This difference has a $\cos \phi$ dependence:

$$d\Delta \sigma_{CH} \equiv d\sigma(e^+p) - d\sigma(e^-p) \sim \cos \phi \times \text{Re} \mathcal{A}.$$  \hspace{1cm} (85)

Here $\mathcal{A}$ is a linear combination of DVCS amplitudes for which the real and imaginary parts are related to corresponding GPDs [189].

5.2.2 HERMES results for DVCS.

The lepton beam-helicity asymmetry and the lepton charge asymmetry in hard electro-production of real photons have been measured for the first time by the HERMES experiment. Data on the beam helicity asymmetry have been published [196]. In this analysis, data taken during 1996 and 1997 with polarised and unpolarised hydrogen targets were used to extract the $(\sin \phi)$ moment of the observed real-photon production cross section. During these data taking periods, the beam positrons were in positive (1996 and 1997) and negative (1997 only) helicity states. The preliminary results on the lepton charge asymmetry based on data taken during 1998 and 2000 with an unpolarised hydrogen target using positron (2000) and electron (1998) beams have been recently presented [197], and are currently being prepared for publication.

The event selection required that apart from the scattered positron (or electron) a photon with energy deposition greater than 0.8 GeV was detected as a single cluster of calorimeter blocks without any additional track in the calorimeter. Besides the usual data quality requirements, fiducial volume cuts and beam polarisation requirements, cuts were also applied to kinematic variables $W^2 > 4 \text{ GeV}^2$, $Q^2 > 1 \text{ GeV}^2$ and $\nu < 24 \text{ GeV}$, selecting the hard scattering region. The average kinematics of the data are $\langle Q^2 \rangle = 2.6 \text{ GeV}^2$, $\langle x \rangle = 0.11$ and $\langle t \rangle = -0.25 \text{ GeV}^2$.

In Fig. 35 the missing mass distribution of the selected events is compared to the results of a Monte Carlo simulation in which photons from the BH process and those from fragmentation processes in deep-inelastic scattering were included. The missing mass resolution of the spectrometer for the DVCS like events is about 0.8 GeV, therefore the exclusivity of the sample cannot be unambiguously determined. There might be some contributions from e. g. the $\Delta$-resonance in the sample. Due to the modest $M_X$
In order to examine the $\phi$ dependence of the interference for case i) the single-spin asymmetry (SSA) defined as

$$ A_{LU} = \frac{1}{\langle P^B \rangle} \frac{[N^+(\phi) - N^-(\phi)]}{[N^+(\phi) + N^-(\phi)]} \tag{86} $$

has been determined. Here $N^+(\phi)$ and $N^-(\phi)$ represent the normalised yields for the respective beam helicity states, and $\langle |P^B| \rangle$ is the average value of the absolute beam polarisation. This SSA for events with missing mass between $-1.5$ GeV and $+1.7$ GeV, i.e. $3\sigma$ below and $1\sigma$ above the nucleon mass is shown in the left panel of Fig. 36.

The comparison of these data with a simple $\sin \phi$-curve demonstrates that the data have the $\phi$-dependence as expected from Eq. (84). The results of the GPD model calculations of Ref. [198] computed at the average kinematics of HERMES are also shown (solid curve). In Fig. 38 (left) the $\langle \sin \phi \rangle$ moment is shown as a function of the missing mass, $M_x$. This moment is defined as:

$$ A_{LU}^{\sin \phi} = \frac{2}{N} \sum_{i=1}^{N} \frac{\sin \phi_i}{\langle P^B \rangle_i}, \tag{87} $$

where the sum runs over all events in a particular bin of missing mass, and $P^B_i$ is the beam polarisation at the time when event $i$ is measured at an angle $\phi_i$. Due to the rather poor missing mass resolution the $M_X$-bins left and right of $M_X = M_{\text{proton}}$ also show non-zero values of $A_{LU}^{\sin \phi}$. Integrated over the missing mass range below 1.7 GeV the analysing power yields:

$$ A_{LU}^{\sin \phi} = -0.23 \pm 0.04(\text{stat.}) \pm 0.03(\text{sys.}). \tag{88} $$

Recently the observation of a non-zero SSA has been confirmed by the CLAS experiment at JLAB [199] at a much lower beam energy of 4.25 GeV. The extracted value of the analysing power of $0.202 \pm 0.028$ is consistent with the HERMES result. The sign is different due to the different beam charge (see Eq. (84)).
Figure 36: HERMES result for the single-spin asymmetry $A_{LU}$ (left) and the charge asymmetry $A_{ch}$ (right) from a proton target as a function of the angle $\phi$. The solid curve in the left panel represents a GPD model calculation [198] computed at the average kinematics of the HERMES experiment. The dashed curve in the left panel is a fit in $\sin \phi$ and the solid curve in the right panel a fit in $\cos \phi$.

As a next step the dependence of $A_{LU}^{sin \phi}$ on $-t$, $x$ and $Q^2$ for the initially analysed data sample was studied. Preliminary results are shown in Fig. 37. Within the given statistical accuracy, no specific dependences can be observed yet.

Also the beam charge asymmetry has been measured for the first time. It is defined as:

$$A_{c}^{\cos \phi} = \frac{A_{c}^{\cos \phi^+} - A_{c}^{\cos \phi^-}}{2}$$

with

$$A_{c}^{\cos \phi^\pm} = \frac{2}{N_{c^{\pm}}} \sum_{i=1}^{N_{c^{\pm}}} \cos \phi_i$$

Preliminary results [197] for the measured charge asymmetry are shown in the right panel of Fig. 36. It exhibits the expected $\cos \phi$ dependence and is significantly different from zero. Fig. 38 shows the preliminary result [197] for the $\cos \phi$ moment as a function of missing mass. In the region below 2 GeV $A_{c}^{\cos \phi}$ is positive with a value of about 0.1.

A drawback of the present HERMES analysis is the relatively poor missing mass resolution for real photons which is dominated by the resolution of the calorimeter. To improve the exclusivity of the
events the HERMES collaboration has decided to install a recoil detector around the storage cell target during the next years. This detector will allow to reconstruct the kinematics of the recoiling particle and to improve the resolution in the squared momentum transfer $t$ by about an order of magnitude. It will also help to reject events in which the real photon is accompanied by a $\Delta$-resonance instead of a proton. The decay products of the $\Delta$ will have a transverse momentum component with respect to the direction of the recoil momentum and thus violate coplanarity with the reaction plane defined by the momenta of the virtual and the real photon.

Fig. 39 shows a cut through a CAD model of the detector. It will consist of an inner part with two layers of Silicon detectors, which are located around the target cell in the machine vacuum, two barrels of Scintillating Fiber Detectors and a Photon Detector. A magnet coil will provide a longitudinal field of 1 T to prevent Möller electrons from hitting the silicon detector and to allow some momentum analysis of the recorded charged particles.

It is planned to take data with unpolarised gas and correspondingly higher target density with this
Recoil Detector for about two years. Fig. 40 shows projections of the statistical accuracy of the beam helicity asymmetry $A_{UL}(\phi)$ (left panel) and of the charge asymmetry $A_{\xi}(\phi)$ (right panel) as a function of azimuthal angle $\phi$ for an integrated luminosity of 2 fb$^{-1}$, together with the results presented above. It is evident that the expected statistical precision will allow a detailed study of the moments of these asymmetries as a function of kinematical variables like $x$ and $t$.

![Figure 40: Projection of the statistical accuracy of $A_{UL}(\phi)$ (left panel) and $A_{\xi}(\phi)$ (right panel) expected with the HERMES Recoil Detector for an integrated luminosity of 2 fb$^{-1}$.

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5.3 Single-spin azimuthal asymmetries in exclusive pion electro-production

The polarised GPDs $\tilde{E}$ and $\tilde{H}$ can be probed through exclusive pseudo-scalar meson electro-production without the need for a polarised target or beam. However, only a quadratic combination of GPDs appears in the unpolarised cross section, and therefore several different observables are needed to disentangle the various contributions [201].

The situation is easier in the case of polarised beams and targets. It has been predicted [202, 129, 203] that for the exclusive production of $\pi^+$ mesons from a transversely polarised target by longitudinally polarised virtual photons the interference between $\tilde{E}$ and $\tilde{H}$ leads to a large target related single-spin asymmetry in the distribution of the angle $\phi$, as defined in the previous section and Fig. 34, the real photon $\gamma$ being replaced by the $\pi^+$ meson.

Although no data have (yet) been collected with transversely polarised targets, some information can already be obtained from a target polarised longitudinally with respect to the beam momentum, since also in that case a small transverse target polarisation component $P_{LT}$ relative to the virtual photon direction appears in addition to the main longitudinal target component:

$$P_{LT} = |P^T| \cdot \sin \theta_\gamma,$$

where $\theta_\gamma$ is the virtual photon emission angle and $\sin \theta_\gamma$ has an average value of 0.16 for the selected HERMES data. Also the virtual photon beam polarisation is longitudinally-transversely mixed instead of being only longitudinal.

In this case the polarised cross section can be written as [187]

$$\sigma_p \sim [P_{LT}^T \sigma_L + P_{LT}^T \sigma_{LT}] \cdot A_{UL}^{\sin \phi} \cdot \sin \phi,$$

where $A_{UL}^{\sin \phi}$ is the sin $\phi$ moment of the polarisation asymmetry for an unpolarised beam (U) and a longitudinally polarised target (L).
While factorisation has been proven for pion electroproduction induced by longitudinally polarised photons [175] there exists at present no proof for electroproduction induced by transversely polarised photons. The interference term $\sigma_{LT}$ can hence not yet be calculated. Quantitative predictions for this term would require a calculation in next-to-leading twist [203].

Figure 41: a) Missing mass distribution for $\pi^+$ (filled circles) and $\pi^-$ (empty circles) electroproduction from a proton target at HERMES. b) Difference between the $\pi^+$ and $\pi^-$ normalised distributions.

HERMES has performed the first measurement of a single-spin asymmetry in the exclusive reaction $e^+ + p \rightarrow e^+ + n + \pi^+ [204]$. Unfortunately a clear identification of the exclusive channel is difficult due to the limited missing mass resolution for charged particle tracks of about 230 MeV, which does not allow to separate directly the exclusive channel from the neighboring (non-exclusive) channels $\pi^+ + \Delta^0$, $\pi^+ + N\pi$, or $\pi^+ + N\pi\pi$. Such non-exclusive channels are also present in the reaction $e^+ + p \rightarrow e^+ + \pi^- + X$, while exclusive $\pi^-$ production on hydrogen with a nucleon in the final state is forbidden by charge conservation. The missing mass distributions for the $\pi^+$ sample (filled circles), and the $\pi^-$ sample (empty circles) are shown in the upper panel of Fig. 41 together with a Monte Carlo prediction for exclusive $\pi^+$ production with arbitrary normalisation (histogram). The $\pi^-$ sample starts at a higher missing mass $M_X$ due to the absence of $\pi^-$ production on hydrogen. The exclusive $\pi^+$ sample has been obtained by subtracting the $\pi^-$ missing mass distributions from the $\pi^+$ distribution after normalising the integral of the $\pi^-$ distribution with respect to the $\pi^+$ one in the range $1.3 \text{ GeV} < M_X < 2.0 \text{ GeV}$ (alternative subtraction schemes yield similar results). The difference between the distributions, presented in the lower panel of the figure shows a clear peak centered at the nucleon mass.

The dependence of the cross section asymmetry for the exclusively produced $\pi^+$ meson on the azimuthal angle $\phi$ is evaluated using

$$A(\phi) = \frac{1}{|P_T|} \frac{N^+(\phi) - N^-(\phi)}{N^+(\phi) + N^-(\phi)}.$$  

where $N$ represents the yield of exclusive $\pi^+$ and the superscript $+(-)$ denotes the target polarisation direction antiparallel (parallel) to the beam momentum. This asymmetry distribution is shown in Fig. 42 versus the angle $\phi$ for all events with $M_X \leq 1.05 \text{ GeV}$ after correcting for the small asymmetry in the background yield. The average values of the kinematic variables are $\langle x \rangle = 0.15$, $\langle Q^2 \rangle = 2.2 \text{ GeV}^2$, and $\langle t \rangle = -0.46 \text{ GeV}^2$. 
Figure 42: The cross section asymmetry $A(\phi)$ averaged over $x, Q^2$, and $t$ for the reaction $e^+ + \bar{p} \rightarrow e'^+ + n + \pi^+$. The data show a large asymmetry with a clear $\sin \phi$ dependence. The curve is a fit of the form $A(\phi) = A_{UL}^{\sin \phi} \cdot \sin \phi$, which yields

$$A_{UL}^{\sin \phi} = -0.18 \pm 0.05 \text{(stat.)} \pm 0.02 \text{(sys)}.$$  

Contributions from $\sin 2\phi$, $\cos \phi$, and $\cos 2\phi$ are compatible with zero.

The dependence of $A_{UL}^{\sin \phi}$ on the kinematic variables $x$, $Q^2$ and $t$ is shown in Fig. 43. The absolute value of the asymmetry shows a clear rise with decreasing $x$. The curves in the left panel represent the upper limits for any asymmetry arising from the transverse target polarisation component $P_T^T$. At low $x$ the measured values are larger than this limit and therefore the asymmetry must arise from the longitudinal component $P_L^L$, independent of any model. It should be noted that it is this low $x$ region where the pion pole is expected [187] to contribute about 70% to the longitudinal cross section. The asymmetry vanishes for small $t$. In the case of a transversely polarised target the second term in

$$\text{Eq. (93) vanishes and so the data expected from future HERMES measurements with a transversely polarised target will give direct access to the GPDs.}$$
In Fig. 44 $A_{UL}^{\text{sin} \phi}$ data are shown for which the exclusivity constraint was released. The sin $\phi$ moment is shown as a function of $z$ and $M_X$ for the $\pi^+$ sample together with preliminary results for $\pi^-$ and $\pi^0$. Obviously something dramatic is happening for $\pi^+$ and $\pi^0$ production when the exclusive region $M_X = M_{\text{nucleon}}$ and $z = 1$ is approached. But unfortunately this region is still very difficult to describe and data from a transversely polarised target, where the effects are expected to be much larger than for a longitudinally polarised target, are urgently needed to experimentally disentangle the asymmetry arising from the two components of the target polarisation.

![Figure 44](image)

**Figure 44:** Preliminary HERMES results for the dependence of $A_{UL}^{\text{sin} \phi}$ for $\pi^+$ production on $^1$H on the missing mass and the variable $z$.

At large $M_X$ (small $z$) there is a small asymmetry which can be related to the transversity distribution, which is discussed in section 6.

### 5.4 Exclusive vector meson production

Above it was stated that: i) Ji's angular momentum sum rule involves the two generalised parton distributions $E(\tilde{z}, \xi, t)$ and $H(\tilde{z}, \xi, t)$ and that ii) exclusive vector meson (VM) production is sensitive to different combinations of these two GPDs. In this respect it is of interest to discuss exclusive VM production in the context of this review, although no spin asymmetries for these reactions have been measured yet, apart from one measurement of a double-spin asymmetry in exclusive $\rho^0$ production [205], which is presented at the end of this section. As the measured (longitudinal) cross sections also provide information on GPDs, some of the main results for exclusive production of the light neutral vector mesons $\rho^0$, $\phi$ and $\omega$ are presented in this section. Details of the analysis can be found in Ref. [206].

The signature of $\rho^0$ and $\phi$ events is simple: a scattered electron together with two oppositely charged hadron tracks. The recoiling target nucleon remains outside of the spectrometer acceptance. A peak corresponding to $\rho^0$ mesons appears in the invariant mass spectrum when the assumption is made that both hadrons are pions, $\phi$ mesons appear with the assumption that they are kaons. The reconstruction of $\omega$ mesons is more complicated. First a $\pi^0$ meson has to be reconstructed from two untracked photon clusters in the calorimeter. Subsequently the invariant mass for the dominant decay channel $\omega \rightarrow \pi^0 \pi^+ \pi^-$, which has a branching ratio of almost 90%, has to be evaluated. Alternatively, for the decay channel $\omega \rightarrow \pi^0 \gamma$, another photon has to be reconstructed. The following kinematical variables are of importance: $p_{\nu}$, $M_\nu$, $M_X$, $\Delta E$, and $t$, which are defined as follows: $p_{\nu}$ is the four-momentum of the VM, $M_\nu = \sqrt{p_{\nu}^2}$ is the reconstructed mass of the VM candidate, $M_X$ is the invariant mass of the undetected final state, $\Delta E = (M_X^2 - M^2)/2M$ is the excitation energy transfered to the target nucleon, and $t = (q - p_{\nu})^2 < 0$ is the squared four-momentum transfer to the target. In the case of an exclusive process, no excitation energy is transfered to the target ($\Delta E \approx 0$). At $t_0$, the maximum (least negative)
value of $t$ allowed for fixed $Q^2$, $\nu$, $M_V$ and $M_X$ the momentum of the final state $VM$ is parallel to the
direction of the incoming photon in the photon-nucleon centre-of-mass system.

Figure 45: Distribution of excitation energy $\Delta E$ transferred to the target for the decay channel $\rho^0 \rightarrow 
\pi^+\pi^-$ based on the HERMES polarised proton data.

The measured $\Delta E$ distribution for the $\rho$-meson sample collected with the polarised hydrogen target
is shown in Fig. 45. Exclusive events are located in the peak around $\Delta E = 0$. The shaded area
$|\Delta E| < 0.6$ GeV indicates the events selected for the analysis of exclusive $\rho^0$ meson production. For
$\Delta E > 2$ GeV the spectrum is dominated by background from non-exclusive processes.

The measured virtual photo-production cross section for exclusive $\rho^0$ mesons [207] is displayed in
Fig. 46 as a function of $W$ for several fixed $Q^2$ values. The HERMES data are seen to fill the gap
between existing measurements at low energy [208, 209, 210], which exhibit a typical fall-off with $W$, and
the data at higher energy [211, 212, 213, 214], which are characterised by a slow rise in $W$. The low
$W$ data can be understood in terms of Reggeon-exchange while for high $W$ Pomeron exchange becomes
dominant. In the intermediate energy domain covered by the HERMES data at $W$ around 5 GeV one
can therefore expect contributions from both types of reaction mechanisms.

Fig. 47 shows the preliminary HERMES results together with results of other experiments for the
cross sections of exclusive $\phi$ (left panel) and $\omega$ (right panel) [216, 215, 210, 217, 218] production on
hydrogen.

The HERMES data cover a so far unexplored $Q^2$ and $W$ range. At low $W$ values the $W$-dependences
of the $\rho$ and the $\omega$ electroproduction cross-sections are very similar, while the $W$ dependence of the
cross section for $\phi$ production looks rather different. There is no rise with decreasing $W$, indicating
that the production mechanism is different.

For a comparison with GPD model predictions the longitudinal cross-section $\sigma_L$ has to be isolated.
This can be determined using the information contained in the angular distribution of the VM decay
products. From these distributions one can extract spin density matrix elements [219] denoted by $r_{ij}^{kl}$.
For the notation see Ref. [220]. Of importance in this context is the matrix element $r_{04}^{04}$. Assuming
s-channel helicity conservation, which implies that the helicity of the virtual photon is preserved by the
produced vector meson, this matrix element is related to the ratio $R$ of longitudinal over transverse
production cross sections

$$ R = \sigma_L / \sigma_T = \frac{r_{04}^{04}}{1 - r_{04}^{04}}. $$  \hspace{1cm} (96)

Preliminary HERMES results for $R$ are shown in Fig. 48 for $\rho^0$ production (left panel) and for $\phi$
production (right panel) as a function of $Q^2$. The data are compared to existing data [221]. While
Figure 46: The cross section for exclusive $\rho^0$ electro-production as a function of $W$ for four different bins in $(Q^2)$. Both the data obtained by HERMES, and those collected by other experiments [208, 209, 210, 211, 212, 213, 214] are displayed.

There is a huge amount of data for $\rho$ production, the corresponding information for $\phi$ production is still rather poor. There is good agreement between the HERMES data and the other data sets in the region of overlap. At high values of $Q^2$ the longitudinal cross-section is seen to dominate over the transverse one.

The longitudinal cross section can be derived from

$$\sigma_L = \frac{R}{1 + \epsilon R} \sigma_{\text{total}},$$

where $\sigma_{\text{total}}$ represents the total measured cross section for exclusive VM production as displayed in Figs. 46 and 47.

The resulting values of $\sigma_L$ are shown in Fig. 49 for $\rho$ production (left panel) and $\phi$ production (right panel) as a function of $W$ for two $Q^2$ bins. In the left panel the high energy muon data from E665 [212] are also shown. The data are compared to the GPD model calculation of Ref. [222]. In the left panel the dotted curves represent the contribution from gluon-exchange, the dashed curves the quark-exchange contribution and the full line is the sum of the two, while the curve in the right panel represents the gluon contribution only. It is important to note that in the case of the $\rho$ vector meson both quark-exchange and gluon-exchange are needed to describe the data, the quark-exchange being dominant in the $W$ range of HERMES, while for the $\phi$ vector meson the gluon-exchange mechanism is sufficient to describe the longitudinal cross section. This might be explained by the facts that the quark content of the $\phi$ is $s\bar{s}$, while the proton contains an almost negligible amount of strange quarks at the $Q^2$ probed by HERMES.

The $\omega$ and the $\rho^0$ vector mesons have similar quark content and their production cross sections are expected to be dominated by the same reaction channels. Indeed their cross section ratio, which is shown in Fig. 50, stays almost constant with $Q^2$ and is within errors compatible with the SU(4) prediction of 1/9. There is an indication of a decrease of the ratio with $W$, which can be explained by additional diagrams for $\omega$ production at low energy. In the case of the $\phi$ the cross section ratio $\sigma_\phi/\sigma_\rho$ is only about one third of the SU(4) expectation 2/9 at HERMES kinematics. It exhibits an increase with both $Q^2$ and $W$ and seems to approach the SU(4) expectation at large values of $Q^2$. The small value of $\sigma_\phi/\sigma_\rho$ is possibly related to the absence of quark-exchange contributions in the $\phi$ case.
Figure 47: Preliminary HERMES data for the cross sections of exclusive \( \phi \) (left) and \( \omega \) (right) electroproduction as a function of \( W \) together with the existing data [210, 215, 216, 217, 218].

Figure 48: Presently available data for the cross section ratio \( R = \sigma_L/\sigma_T \) obtained from the decay angular distributions of diffractive \( \rho^0 \) production (left). Preliminary HERMES results together with other data for exclusive \( \phi \) electroproduction (right) [221].
Figure 49: The derived longitudinal cross section for exclusive $\rho^0$ (left) and $\phi$ (right) electroproduction together with predictions from GPD-based calculations, and some existing data.

Figure 50: Preliminary HERMES results for the cross section ratio for exclusive $\omega$ and $\rho^0$ electroproduction (left) and for exclusive $\phi$ and $\rho^0$ production (right) as a function of $Q^2$. Existing data are also shown.
5.5 Double-Spin Asymmetry in the Cross Section for Exclusive $\rho^0$ Electroproduction

With its polarised beam and target HERMES has the ability to also investigate double-spin asymmetries in vector meson production. Here the results for a double-spin asymmetry obtained for exclusive $\rho^0$ production [161, 205] are presented. From the combined 1996 and 1997 data on the polarised hydrogen target about 2800 events remained after the following selection criteria had been applied: $0.62 \text{ GeV} < M_{\pi\pi} < 0.92 \text{ GeV}$, $-t' < 0.4 \text{ GeV}^2$ and $|\Delta E| < 0.6 \text{ GeV}$. The average values of the relevant kinematic variables of this data sample are: $\langle W \rangle = 4.9 \text{ GeV}$, $\langle Q^2 \rangle = 1.7 \text{ GeV}^2$, $\langle -t \rangle = 0.1 \text{ GeV}^2$, $\langle x \rangle = 0.07$. The asymmetry $A_1^p$ has been determined from the asymmetry $A_{\|}$ using the positivity limit $A_1^p = \sqrt{R(Q^2)}$, where a fit to the data for $R$ obtained from the decay angular distributions [219] has been used for $R$. For the present data sample this fit yields an average value of $R = 0.62$. The evidence for a positive double-spin asymmetry is shown in Fig. 51, where $A_1^p$ is plotted versus $Q^2$, $W$, $x$, and $-t'$.

![Figure 51: HERMES results for the double-spin asymmetry in $\rho^0$ production from a polarised hydrogen target as a function of the kinematical variables $Q^2$, $W$, $x$, and $-t'$ [161, 205].](image)

Error bars denote the statistical uncertainties and the band at the bottom of the plots indicates the systematic uncertainties which are dominated by the contributions from the beam and target polarisation measurements. Within the statistical accuracy of the data no significant dependence on any of the kinematical variables or angles can be seen. Averaged over the kinematic acceptance the measured asymmetry has a value

$$\langle A_1^p \rangle = 0.24 \pm 0.11(\text{stat.}) \pm 0.02(\text{syst.}).$$

(98)

The result is compatible with a prediction which is based on the description of diffractive exclusive $\rho$ leptoproduction and inclusive deep-inelastic scattering by the generalised vector meson dominance model [223]. It relates the asymmetry for exclusive $\rho$ leptoproduction $A_1^p$ to the inclusive asymmetry $A_1$. From the predictions presented in Ref. [223] a ratio of about 2 can be estimated for the ratio $A_\rho/A_1$ at the HERMES beam energy, which is in good agreement with the experimental value $1.9 \pm 0.8$. Here the non-vanishing asymmetry indicates the presence of production amplitudes corresponding to unnatural parity exchange in the $t$-channel. This can be considered to originate from diquark exchange which agrees with our previous conclusion that longitudinal $\rho^0$ production at HERMES energies is dominated by quark exchange. The measured non-zero asymmetry is in conflict with a preliminary result of a similar measurement by the SMC collaboration [224] at comparable values of $Q^2$ but at three times higher values of $W$, which is consistent with zero. This apparent discrepancy can, however, be understood by two facts: (i) at higher values of $W$ the production process might be dominated by gluon exchange rather than quark exchange; and (ii) at the smaller values of $x$ probed by SMC the inclusive asymmetry is much smaller.

Since the proof for factorisation of the hard and soft amplitudes exists only for longitudinal photons and does not apply to the production of mesons from transverse photons [175, 225], a clear interpretation of the observed double-spin asymmetry in $\rho^0$ production within the framework of perturbative QCD and generalised parton distributions does not exist yet and would require substantial theoretical progress.
Single Spin asymmetries in semi-inclusive deep-inelastic electro-production of pions - transversity

More than twenty years after the first deep-inelastic measurements became available it was realised [226, 227, 228, 229, 230] that a complete description of the partonic structure of the nucleon in leading twist requires in addition to the unpolarised quark distributions \( q(z, Q^2) \) and the quark helicity distributions \( \Delta q(x, Q^2) \), a third quark distribution, named transversity \( \delta q(x, Q^2) \). The corresponding structure function is denoted as \( h_1(x, Q^2) \). (In fact transversity was first identified about 10 years earlier by Ralston and Soper [231], when discussing polarised Drell-Yan reactions, but then was somehow forgotten again).

Pedagogical introductions into the topic are given for instance in Refs. [232, 233, 234], while the present status of the field is summarised in a recent very comprehensive review [235]. Here we only summarise some of the main arguments.

Physically, one may think of transversity \( \delta q(z) \) as the distribution of quarks with spin parallel to the nucleon spin minus the distribution of quarks polarised anti-parallel to the nucleon spin for a transversely polarised nucleon, in analogy to the helicity distribution \( A_q(z) \) for a longitudinally polarised nucleon. For quarks moving relativistically in a nucleon at infinite momentum \( \delta q(x) \) and \( \Delta q(x) \), which should be identical in the non-relativistic quark model, are different since rotations and boost do not commute in the relativistic case and therefore the distributions cannot be converted into each other by a series of boosts and rotations, as in the non-relativistic case.

The transversity \( \delta q(x) \) can be related to a helicity amplitude \( A_{hH,h'H'} = A_{+-,-+} (h, H, h', H') \) the helicities of lepton \((h, h')\) and nucleon \((H, H')\) in the initial and final states with absolute value \(1/2\) and signs + and -, where a helicity flip occurs for both the lepton and the nucleon, while \( q(x)(\Delta q(x)) \) is related to the sum (difference) of the only two independent helicity amplitudes \( A_{+-,++} \) and \( A_{+-,+-} \), allowed by angular momentum, parity and time reversal invariance, where no helicity flip occurs. Therefore, the transversity distribution \( \delta q(x) \) is (as opposed to \( q(x) \) and \( \Delta q(x) \)) a chiral-odd object.

Since electroweak and hard QCD processes preserve helicity, \( \delta q(x) \) is not accessible in inclusive deep-inelastic scattering, which is the main reason why it escaped experimental observation until now and is still unmeasured.

Transversity can only be measured in conjunction with another chiral-odd object like e.g. a chiral-odd anti-quark transversity distribution in transversely polarised Drell-Yan production of lepton pairs, or a chiral-odd fragmentation function in the current fragmentation region of semi-inclusive deep-inelastic scattering. We come back to the second option below, and first summarise the expected main properties of transversity.

- Due to helicity conservation there is no transverse analogue of the longitudinal polarised gluon distribution \( \Delta g(x) \) in a spin-1/2 target. Hence \( \delta g(x) \) does not mix with gluons under evolution. Consequently, \( \delta q(x) \) and \( \Delta q(x) \) are predicted to evolve very differently with \( Q^2 \), the transverse quarks just loose momentum by gluon emission, but the distributions do not get contributions at low \( x \) from gluon annihilation into pairs of transversely polarised quarks and anti-quarks.

- The integral of the transversity distribution \( \delta \Sigma (Q^2) = \sum_q \int_0^1 dx \left( \delta q(x, Q^2) - \delta \bar{q}(x, Q^2) \right) \) is called 'tensor charge' due to the tensor character of the corresponding local operator, should be much closer to the relativistic quark model value of about 0.6 than \( \Delta \Sigma (Q^2) \). Indeed calculations within various models (see Table 6 in Ref. [235] and corresponding references) consistently arrive at values around 0.6. Note, however, that the tensor charge is not conserved but will decrease with \( Q^2 \) due to the evolution of \( \delta q(x, Q^2) \).

- Finally there are (in leading order QCD) some important inequalities relating the three twist-2 distribution functions: \( |\delta q(x, Q^2)| \leq q(x, Q^2) \) and the Soffer inequality [236] \( |2\delta q(x, Q^2)| \leq q(x, Q^2) + \Delta q(x, Q^2) \).
One of the possibilities to study transversity [235, 233] is via the azimuthal angular asymmetry in the distribution of hadrons produced in semi-inclusive deep-inelastic scattering via the chiral-odd fragmentation function $H_{1}^{x}$, the so-called Collins fragmentation function [237, 238]. The Collins fragmentation function will arise if there are interferences of different reaction channels, i.e. final state interactions, generating a nontrivial phase in the DIS amplitude [239, 240, 241]. If $H_{1}^{x}$ is not too small, which is suggested from DELPHI data on azimuthal correlations between particles produced from opposite jets in $Z^{0}$ decay [242] then the measurement of such an asymmetry will enable the extraction of the transversity distribution.

For an unpolarised beam (U) and a transversely polarised target (T) the corresponding azimuthal asymmetry for a meson is given by

$$A_{UT}^{s} = a \cdot \frac{2(1 - y)}{1 + (1 - y)^2} \cdot \sum_{q} e_{q}^{2} \delta q(x) H_{1,q}^{l,m}(z),$$

$$= C_{m} \cdot \sum_{q} e_{q}^{2} \delta q(x) H_{1,q}^{l,m}(z),$$

where $a$ is a coefficient depending on the assumption made for the dependence of the transversity distribution and the fragmentation function on the transverse momenta of the quarks in the nucleon and of the outgoing meson. Note that the determination of the transversity distribution for the various quark flavours is much more complicated than the determination of the spin-dependent parton distributions discussed in Sect. 4.4, since both the parton distributions $\delta q(x)$ and the fragmentation functions $H_{1,q}^{l,m}(z)$ are unknown. Measurements of azimuthal asymmetries for different kinds of mesons and for hydrogen and deuterium targets will be required to entangle them.

As in the case of exclusive electro-production, discussed in the preceding section, an azimuthal asymmetry can also arise in the case of a longitudinally polarised target from the transverse component $P_{L}$ of the target polarisation with respect to the direction of the virtual photon. Here $P_{L}$ lies in the scattering plane and the angle $\phi$ is then again the angle between the scattering plane and the hadron production plane. In the case of a longitudinal polarised target the sin $\phi$-weighted asymmetry thus contains additional leading and higher twist terms

$$A_{UL}^{s} = \frac{2(1 - y)}{1 + (1 - y)^2} \cdot \sum_{q} e_{q}^{2} \delta q(x) H_{1,q}^{l,m}(z)$$

$$+ \frac{P_{L} M}{Q} \frac{2(1 - y)}{1 + (1 - y) \left( \sum_{q} e_{q}^{2} \delta q(x) H_{1,q}^{l,m}(z) + \ldots \right)},$$

where $\delta q_{L}(x)$ (or the corresponding structure function $h_{1,q}^{L}(x)$) is one of the additional distribution functions discussed in detail in Ref. [240]. Additional terms contain contributions of $h_{1,L}^{L}(x)$ and $H_{1,q}^{L}$ with interaction dependent twist–3 functions $h_{1}$ and $H_{1}$. Note that with our definition of the angle $\phi$ the transverse contribution appears with a negative sign [243, 244] and that the kinematic factor in front of the longitudinal term is about one to two orders of magnitude larger than the one in front of the transverse term.

HERMES has observed such an azimuthal asymmetry in the semi-inclusive electro-production of charged [245] and neutral [246] pions on a longitudinally polarised target. The identification of charged pions was discussed above. Neutral pion identification was provided by the detection of two photon clusters in the electromagnetic calorimeter originating from $\pi^{0}$ decay, each with a minimum energy deposition of 1.0 GeV and without a corresponding charged track. Fig. 52 (left) shows the reconstructed two-photon invariant mass distribution. Neutral pions were selected within the invariant mass range of $0.10 \text{ GeV}^{2} < m_{\gamma\gamma} < 0.17 \text{ GeV}^{2}$, where background contributions from uncorrelated photons typically amount to 20%.

Fig. 52 (right) shows the measured azimuthal asymmetry from the measurements with the longitudinally polarised proton target in 1996/97 as a function of the azimuthal angle $\phi$ for $\pi^{+}$, $\pi^{-}$ and $\pi^{0}$ production. A clear sinusoidal behaviour for $\pi^{+}$ and $\pi^{0}$ production is observed with very similar amplitudes $A_{UL}^{s}(\pi^{+}) = 0.020 \pm 0.004 \text{(stat.)} \pm 0.003 \text{(syst.)}$ and $A_{UL}^{s}(\pi^{0}) = 0.019 \pm 0.007 \text{(stat.)} \pm 0.003 \text{(syst.)}$, respectively.
The asymmetry for $\pi^{-}$ production is measured to be compatible with zero within errors. This can be understood from the fact that the up quark distribution dominates in the nucleon and appears with a four times larger charge factor in the cross section than that of the down and strange quarks. Moreover, the fragmentation from up quarks into $\pi^{+} = (u\bar{d})$ is much more likely than from up quarks into $\pi^{-} = (d\bar{u})$ (unfavoured fragmentation function), while for the $\pi^{0}$ both up and down quark distributions contribute with favoured fragmentation functions. The dependence of the amplitude $A_{UL}^{\sin \phi}$ on the variables $z$, $x$, and $p_{T}$ is shown in Fig. 53 after averaging over the other two kinematic variables. If one assumes the factorisation ansatz in terms of the transversity distribution $\delta q(x)$ and the Collins fragmentation function $H_{i}^{L}(z)$ to be correct then the non-zero asymmetries for $\pi^{+}$ and $\pi^{0}$ production provide evidence in support of their existence and imply a substantial magnitude for both of them. The positive value of the asymmetry indicates that the longitudinal term in Eq. 100 makes the dominant
contribution to $A_{UU}^{\sin \phi}$. The dependence on $p_\perp$ can be related to the dominant kinematic role of intrinsic transverse momentum of quarks, if $p_\perp$ remains below the typical hadronic scale of around 1 GeV. Also shown in Fig. 53 are predictions of a model calculation for the $\pi^+$ and $\pi^0$ single-spin azimuthal asymmetry [247, 248], using isospin and charge conjugation invariance [249]. In this model calculation, however, the longitudinal contribution and the transversal enter with the same sign. Altogether these data provide information and constraints also for other phenomenological approaches [235] like the chiral-quark soliton model [244, 250], the quark-diquark model or a pQCD model [251].

Preliminary HERMES results on such single-spin asymmetries are now also available from the measurements with a longitudinally polarised deuterium target in 1998/2000 as well for pions and kaons, which were identified with the RICH [252]. Examples are shown in Fig. 54.

![Figure 54: Preliminary HERMES results on the analysing power $A_{UL}^{\sin \phi}$ for charged pion and kaon production from a deuterium target.](image)

The $\pi^+$ and $\pi^0$ asymmetries for the deuteron are generally smaller than those for the proton and the $\pi^-$ asymmetries are now different from zero, compatible with the expectations from the quark model and also with recent theoretical predictions [250, 253].

The interpretation of the data in terms of the transversity distribution and the Collins fragmentation function is presently under dispute. Recently, it has been argued [254] that a single-spin azimuthal asymmetry can arise from a final state interaction mediated by gluon exchange between the outgoing quark and the target spectator system at leading twist in QCD and that in this case the single-spin azimuthal asymmetry does not factorise, as the product of quark distribution and fragmentation function does. In particular the claim is that this asymmetry need not be related to the transversity distribution.

Indeed such an angular asymmetry could also be produced by an unknown T-odd distribution function $f_{1T}^q$ [240], first postulated by Sivers [255, 256] together with the familiar fragmentation function $D(z)$. As discussed in [243, 257] future measurements with a transversely polarised target will allow to distinguish between the Collins mechanism and the Sivers mechanism [255, 256, 258, 259, 260]. The Collins mechanism will cause a $\sin(\phi + \phi_S)$ moment proportional to $h_1(x)H_T^1(z)$, while the Sivers mechanism will cause a $\sin(\phi - \phi_S)$ moment proportional to $f_{1T}^q(x)D(z)$. Here $\phi_S$ is the angle between the electron scattering plane and the target polarisation plane introduced in Sect. 3.1. In this context one should also mention the 'rotating constituent' model by Meng [186]. In this very qualitative model [261] the orbital motion of the quarks in the proton around the polarisation axis produces an azimuthal angular distribution of leading hadrons with maxima in the plane perpendicular to the target polarisation plane.

It is obvious that much more detailed experimental data are needed to investigate the kinematical dependences of these asymmetries and to confront the extracted transversity distributions and fragmentation functions with theoretical predictions.

HERMES recently started to take data with a transversely polarised hydrogen target. The effects are expected to be much larger than for a longitudinally polarised target and the azimuthal asymmetries
Figure 55: Expectations for the accuracies of the weighted transverse asymmetry $A_T^{+\pi}$ for semi-inclusive $\pi^+$ production (left panel) and the transversity distribution $\delta u(x)$ for up quarks (middle panel) as a function of $x$, and the ratio of the fragmentation functions $H_{1u}^{(1),+\pi}(z)$ and $D_u^{+\pi}(z)$ for 2 years of HERMES data taking with the transversely polarised hydrogen target [262]. Note that the analysing power $A_T^{+\pi}$ is different from Eq. (99) as it is weighted by $p_L/(z m_\pi)$, where $p_L$ is the transverse momentum of the pion and $m_\pi$ its mass. The superscript (1) appearing on the function $H_{1u}^{(1),+\pi}$ denotes a $p_L$-weighted integral over transverse momentum of the pion.

are free from longitudinal contributions such as the second term in Eq. 100. Fig. 55 shows predictions of the estimated accuracies which can be obtained within 2 years of data taking [262]. The $Q^2$ dependence of $\delta q(x,Q^2)$ can be investigated, when also data from the COMPASS experiment at six times higher beam energy will be available. It is worth noting that the detailed study of transversity is the main objective for the proposed fixed target experiment TFSI.A-N [263].
7 Spin Phenomena in Lambda Electroproduction

The \( \Lambda \) hyperon is an especially interesting particle in spin physics. It is the lightest strange baryon \( (M_\Lambda = 1115.684 \text{ MeV}) \) and thus can only decay weakly \( (\tau = 2.632 \times 10^{-16} \text{ s}) \), the dominant decay channels being \( \Lambda \rightarrow pn^- \) (63.9%) and \( \Lambda \rightarrow n\pi^0 \) (35.8%). Since in the HERMES experiment \( \Lambda \) particles are only detected via the charged decay channel, we will only discuss this channel in the present section. As the weak decay is parity violating, the proton is being emitted preferentially into the direction of the \( \Lambda \) spin. The angular distribution of the decay nucleon in the \( \Lambda \) rest frame has the form

\[
\frac{dN}{d\Omega_p} \propto 1 + \alpha \vec{F}_\Lambda \cdot \hat{k}^p.
\]

Here, \( \hat{k}^p \) is the unit vector along the proton momentum in the \( \Lambda \) rest frame, \( \vec{F}_\Lambda \) is the polarisation of the \( \Lambda \) and \( \alpha = 0.642 \pm 0.013 \) [264] is the asymmetry parameter of the weak decay. Information about the \( \Lambda \) polarisation can therefore be directly obtained from the angular distribution of the decay particles.

In the naive QPM all the spin of the \( \Lambda \) is carried by the strange quark \( (\Delta s^\Lambda = 1) \), the up and down quark forming a spin singlet, \( (\Delta u^\Lambda = 0) \) and \( (\Delta d^\Lambda = 0) \). On the other hand the \( \Lambda \) is a member of the hyperon octet and based on SU(3) flavour symmetry all the quark contributions to the spin of the \( \Lambda \) are related to those of all other members in the octet including the proton and the neutron. The nucleon data tell us that only a small fraction of the nucleon spin originates from the spins of its quarks, \( \Delta \Sigma \approx 0.2 \), in contrast to the naive QPM prediction, \( \Delta \Sigma = 1 \). As a consequence, also the up and down quarks should be expected to give sizeable contributions to the spin of the \( \Lambda \) hyperon [265].

The assumption \( \Delta \Sigma^\Lambda = 0.18 \), together with the values for the weak decay constants \( F \) and \( D \) and the weak axial charge \( g_A = g_A/g_V \) resulted in the estimate: \( \Delta u^\Lambda = -0.2, \Delta d^\Lambda = -0.2 \) and \( \Delta s^\Lambda = 0.58 \) (from Ref. [265]). Lattice calculations of the spin structure of the \( \Lambda \) hyperon [266], however, predict very small values of \( \Delta u^\Lambda \) and \( \Delta d^\Lambda \) compatible with zero.

7.1 Longitudinal spin transfer in \( \Lambda \) electroproduction

Information about the spin structure of the \( \Lambda \) hyperon can be obtained in spin transfer experiments. First a longitudinally polarised quark is produced from a polarised intermediate photon (or \( W^\pm \) boson) in polarised deep inelastic electron or muon (neutrino) nucleon scattering (as explained in section 3.3) or from \( Z^0 \) decay on the \( Z^0 \) pole. This longitudinally polarised quark fragments into a \( \Lambda \) particle. If the spin is preserved in the fragmentation process, independent of the quark flavour, then the \( \Lambda \) will also be longitudinally polarised and the degree of longitudinal \( \Lambda \) polarisation can be determined from the angular decay distribution. Obviously in such experiments the measured polarisation does depend on both the \( \Lambda \) spin structure and the spin transfer in the fragmentation process.

For an unpolarised target and a polarised electron beam the polarisation of the produced \( \Lambda \) in the current fragmentation region is [267, 268]:

\[
P^\Lambda = P^B D(y) \sum_q e_q^2 \frac{1}{2} q(x, Q^2) G_{1,q}^\Lambda (z, Q^2),
\]

where \( P^\Lambda \) and \( P^B \) are the polarisations of the \( \Lambda \) and the beam, \( D(y) \) is the depolarisation factor of the virtual photon (Eq. (20)), \( q(x, Q^2) \) is the unpolarised quark distribution in the target nucleon, where \( q \) runs over the quark flavours up, down and strange, and \( D_{1,q}^\Lambda (z, Q^2) \) and \( G_{1,q}^\Lambda (z, Q^2) \) are the unpolarised and polarised fragmentation functions which are a measure of the probability that a quark of type \( q \) fragments into an unpolarised \( \Lambda \), respectively a polarised one.

Since the charge factor for the up quark is four times bigger than the one for down and strange quarks, and the unpolarised up quark distribution is about twice as large as the unpolarised down quark distribution the deep inelastic scattering process is dominated by the contribution from the up quarks. Moreover, due to isospin symmetry the spin transfer from the up and down quarks to the \( \Lambda \) are expected to be equal. Therefore Eq. (102) reduces approximately to

\[
P^\Lambda = P^B D(y) \frac{G_{1,u}^\Lambda (z, Q^2)}{D_{u}^\Lambda (z, Q^2)}.
\]
From this expression we obtain an approximate relation for the so called spin transfer $S^\Lambda$ from the photon to the $\Lambda$

$$S^\Lambda = \frac{P^\Lambda}{P^B D(y)} \approx \frac{G_{1u}^\Lambda(z, Q^2)}{G_{1u}^\Lambda(z, Q^2)}.$$  

(104)

Consequently, $\Lambda$ electroproduction in the current fragmentation region is mostly sensitive to the ratio $G_{1u}^\Lambda(z, Q^2)/D_u^\Lambda(z, Q^2)$ i.e. the spin transfer from the up quark to the $\Lambda$.

From SU(3) flavour symmetry [265, 267] one would expect the average spin transfer $S^\Lambda$ to be approximately equal to $-0.2$. From the naive constituent quark model one would expect a small positive contribution to the spin transfer from the small fraction of $s$ quarks in the nucleon sea contributing fully to the spin of the $\Lambda$, but this will occur only at small values of $x$. Recent calculations [269] based on SU(3) flavour symmetry relations and experimental spin dependent quark distributions predict a positive polarisation of up and down quarks in the $\Lambda$ at high values of $x$ and increasing slightly with $z$. The same authors have previously also calculated the spin transfer in a pQCD based analysis and in a quark-diquark model [270], obtaining a positive spin transfer increasing substantially with increasing $z$.

From all data collected in the years 1996-2000 $\Lambda$ hyperons have been selected through their decay into a proton and $\pi^-$ meson. At least three reconstructed tracks were required: a scattered lepton track in coincidence with two hadron candidate tracks with opposite charge. To ensure that the event occurred in the target gas, the longitudinal vertex position was required to be within the total length of the target cell. The positron interaction vertex and the $\Lambda$ decay vertex were required to be separated by more than 10 cm to eliminate background hadrons originating from the primary vertex. The $\Lambda$ hyperon was identified by reconstructing the invariant mass of the $p\pi^-$ system. The following cuts were applied on the scattered lepton kinematics: $Q^2 > 1 \text{GeV}^2$, $y < 0.85$ and $W > 2 \text{GeV}$. To ensure that the $\Lambda$ hyperons primarily originate from the current fragmentation region, a positive value of $x_F \times 2\pi L_c / W$ was required.

The acceptance of the HERMES spectrometer for the reconstruction of the $\Lambda$ hyperons is limited and strongly depends on $\cos \theta_{pL}$, where $\theta_{pL}$ is the angle between the proton momentum and the $\Lambda$ spin quantisation axis. The HERA lepton beam is always longitudinally polarised at the location of the HERMES experiment and data have been taken with both beam helicities. To minimise acceptance effects, the spin transfer to the $\Lambda$ has been determined by combining the two data sets measured with opposite helicity in such a way that the luminosity-weighted average beam polarisation for the selected data sample is zero. In this case the spin transfer to the $\Lambda$ can be determined from the forward-backward asymmetry of the angular decay distribution [271, 272, 273]:

$$S^\Lambda = \frac{1}{\alpha((P)^2)} \cdot \sum_{i=1}^{N_A} \frac{P_i^B \cos \theta_{pL}^i}{\sum_{i=1}^{N_A} D(y_i) \cos^2 \theta_{pL}^i},$$  

(105)

where the sums are over the $\Lambda$ events and $((P)^2)$ is the luminosity-weighted average of the square of the beam polarisation.

Results from the 1996/97 data, based on about 2200 $\Lambda$ events, have already been published [274]. For an average fractional energy transfer $\langle z \rangle = 0.45$, the longitudinal spin transfer along the $\Lambda$ momentum direction has been found to be $S^\Lambda = 0.11 \pm 0.17\text{(stat)} \pm 0.03\text{(sys)}$.

The total 'helicity balanced' data sample obtained in the years 1996-2000 contains about 9300 $\Lambda$ events in the current fragmentation region ($x_F > 0$). From these data a preliminary value for the average spin transfer of $0.04 \pm 0.09\text{(stat)} \pm 0.03\text{(sys)}$ has been determined. The increased statistics permits to split the data into several bins, with similar statistics in each bin. In Fig. 56 this preliminary result is shown as a function of $z$. The experimental results are, within errors, in qualitative agreement with the expectations of SU(3) flavour symmetry [265, 267] and the naive constituent quark model, but also with the prediction of Ref. [269]. The HERMES data do, however, not agree with a spin transfer increasing with $z$, predicted by some of the models. Obviously, substantially higher statistics data at $z$ close to 1 are needed, and actually of crucial importance for a precise test of the various models for $\Lambda$ spin structure and the spin transfer mechanism.

In Fig. 57 the preliminary HERMES data for the spin transfer as a function of $x_F$ are presented together with recent results from the NOMAD neutrino experiment [275], older data from WA59 [276]
and from the high energy muon experiment E665 [277] at FNAL. The HERMES data extends much further into positive $x_F$ where no other data are available so far. There is nice agreement with the NOMAD data in the region of overlap. It should be noted that the mechanism for spin transfer at $x_F < 0$ could be entirely different [278, 279].

As briefly mentioned in the description of the detector components the 'LambdaWheels' detector [280, 281, 282] has been installed in 2001 for the second data taking period of HERMES 2002–2006. This detector enlarges the geometrical acceptance for $\Lambda$ detection substantially and also extends the kinematic range in $x_F$ to small and negative values (see Fig. 58), such that now also the target fragmentation region can be accessed. This is of special interest, as it has been shown [278, 279] that the longitudinal polarisation of $\Lambda$ particles in the target fragmentation region provides information on the still very poorly known polarisation of the strange sea-quarks in the nucleon.

7.2 Transverse $\Lambda$ polarisation in quasi-real photo-production

Transverse polarisation of $\Lambda$ hyperons has been studied extensively in hadron scattering experiments [283, 284, 285]. In almost all experiments a negative polarisation with respect to the normal of the production plane $\hat{n} \equiv \vec{p}^B \times \vec{p}^\Lambda$ has been observed, increasing almost linearly with transverse momentum.
p_T of the Λ up to a value of about 1 GeV, where a plateau is reached. The observed polarisation also rises linearly with x_F. The only exception came from measurements in K^- p reactions, where the incoming beam contains (valence) strange quarks. Here, a positive polarisation has been observed. Possible mechanisms for the origin of this transverse polarisation have for example been reviewed in Refs. [283, 286]. As yet, no model is able to account for all aspects observed in hyperon and anti-hyperon polarisation.

While transverse Λ polarisation in hadron-hadron interactions at high energies has been studied extensively, very little information exists about this effect in photo- and electroproduction. Apart from some measurements of exclusive Λ production at low energies [287, 288], this phenomenon was only investigated with rather poor statistical accuracy 20 years ago at CERN [289], for incident tagged photon energies between 20 and 75 GeV, and at SLAC [290], using a 20 GeV photon beam. CERN measured an average transverse polarisation of (6 ± 4)% and SLAC of (9 ± 7)%.

HERMES has measured transverse polarisation of Λ and Λ hyperons in quasi-real photoproduction. The scattered electron was not required in the trigger and therefore the data are dominated by events with very low Q^2 values. Hence, the kinematics of the quasi-real photons are not known and only kinematic variables related to the eN-system are available. The average kinematics of the Λ and Λ have been estimated with a Monte Carlo Simulation. An average virtual photon energy of (ν) = 16 GeV has been obtained, comparable to the kinematics of the CERN and SLAC experiments. With similar cuts and particle identification criteria, but of course without the requirements in Q^2, W and y, a data sample of about 386,000 Λ and 72,000 Λ events has been selected in a ±2.5σ window around the Λ (Λ) mass in the pπ^- (pπ^+) invariant mass distribution. A background-subtraction procedure was applied taking into account the false asymmetry of the background, which enters the systematic error of the measurement. Due to the up/down mirror symmetry of the HERMES spectrometer, many acceptance effects cancel and the transverse Λ polarisation can be determined by evaluating

\[ P_\Lambda = \frac{1}{2} \frac{\langle \cos \theta_p \rangle}{\alpha \langle \cos^2 \theta_p \rangle}, \]

where θ_p denotes the angle of proton emission relative to the n axis in the Λ rest frame.

The preliminary result for the measured transverse polarisations are

\[ P_\Lambda = (+5.5 \pm 0.6 \text{ (stat)} \pm 1.6 \text{ (sys)}) \% \quad \text{and} \quad P_\bar{\Lambda} = (-4.3 \pm 1.3 \text{ (stat)} \pm 1.2 \text{ (sys)}) \%. \]

The signs of the measured polarisations are a surprise. The transverse Lambda polarisation is positive, in contrast to almost all other measurements of inclusive Λ production in hadron-nucleon interactions.

Figure 58: Accepted number of events (in arbitrary units) as a function of x_F for the HERMES detector without (data points) and with (histogram) the new Lambda Wheels detector, which is scheduled to be used at HERMES in the years 2002-2006.
apart from those with a $K^-$ beam. The data favour a negative $\Lambda$ polarisation in contrast to the zero value consistently measured in other experiments apart from $p\bar{p} \rightarrow \Lambda\Lambda$.

![Figure 59: Preliminary HERMES results for the transverse $\Lambda$ polarisation in quasi-real photoproduction as a function of the variable $(E_\Lambda + p_z)/(E_b + p_b)$ (left panel), and as a function of $p_T$ for two regions of $\cos \theta_\Lambda$ (right panel), where $\theta_\Lambda$ is the angle between the beam direction and the produced $\Lambda$ in the electron-nucleon center-of-mass frame.](image)

In Fig. 59 (left) the preliminary HERMES results for the transverse $\Lambda$ polarisation are shown as a function of the variable $\xi = (E_\Lambda + p_z)/(E_b + p_b)$, where $p_z$ is the longitudinal momentum of the $\Lambda$ with energy $E_\Lambda$ and $E_b, p_b$ are the energy and momentum of the beam. It can be shown, that for HERMES kinematics positive values of $x_F$ correspond to $\xi$ larger than about 0.25. Consequently for $\xi$ larger than this value the events originate dominantly from the current fragmentation region, the smaller $\xi$, the larger the contribution from the target fragmentation region. Another variable to approximately separate events from the current and target fragmentation region is $\cos \theta_\Lambda$, where $\theta_\Lambda$ is the angle between the beam direction and the produced $\Lambda$ in the $eN$ center-of-mass frame. The (somewhat arbitrary) value of $\cos \theta_\Lambda = 0.6$ again separates roughly events from the current and the target fragmentation region. Fig. 59 (right panel) shows the preliminary HERMES result as a function of transverse $\Lambda$ momentum $p_T$ for the two data samples above and below this cut.

Transverse $\Lambda$ polarisation in electroproduction at high $Q^2$ might originate from the T-odd fragmentation function $D_{1T}$, one of the eight fragmentation functions in semi-inclusive deep-inelastic scattering identified in [240]. The possible origin in quasi-real photoproduction, however, is unclear. The large positive value of the transverse $\Lambda$ polarisation is surprising in the light of the negative values observed in almost all other reactions apart from those with $K^-$ beams. The result might thus indicate that the $\gamma \rightarrow s\bar{s}$ hadronic component of the exchanged photon might be the dominant source for this positive transverse polarisation. But at present these are only speculations. Much more data for $\Lambda$ and $\bar{\Lambda}$ polarisation in both the current and the target fragmentation region are needed to study in detail the kinematical dependences of this polarisation. This is part of the HERMES program in the coming years when the acceptance for $\Lambda$ particles will be substantially enlarged by the new 'LambdaWheels' detector.

It has been proposed [291] to use transverse $\Lambda$ polarisation measured in $\Lambda$ electroproduction from a transversely polarised proton target as a 'polarimeter' for the transversity distribution $\delta q(x)$. Since substantial transverse $\Lambda$ polarisation is already observed for an unpolarised target it will, however, be rather difficult to disentangle the different sources for the transverse polarisation measured in this case. Furthermore, the measured longitudinal spin-transfer from the struck quark is small. This indicates that the transverse spin-transfer will also be small and that thus transverse $\Lambda$ polarisation will hardly serve as a useful tool for future transversity measurements.
8 Outlook

In the richest sense of the word, the HERMES experiment is a facility to study reactions induced by quasi-real and virtual photons from polarised atomic targets and from unpolarised nuclear targets. In the first phase the experiment has yielded a lot of high precision results on polarised deep-inelastic scattering and has confirmed and extended our understanding of the contribution of quark spins to the spin of the nucleon. The experiment has yielded a number of surprises: a first indication of a positively spin-dependent gluon distribution from pairs of hadrons with high transverse momentum, an experimental indication for effects of transversity in the nucleon, azimuthal single-spin asymmetries in hard exclusive electro-production of real photons and pions which can be related to generalised parton distributions, and a positive transverse polarisation of Λ particles produced in quasi-real photo-production. Many more results have been presented in this review, like the measurement of a flavour asymmetry in the light quark sea and cross-sections for hard exclusive vector meson production. In addition several nuclear effects in deep-inelastic hadron production from nuclear targets have been investigated [292, 293] but were not discussed in the present review. Many more results can be expected from ongoing analyses based on data collected in this first phase of HERMES (1995-2000).

In the second phase of HERMES in the years 2002-2006 a transversely polarised proton target will be used for about two years to perform a first measurement of the transversity distribution and transverse fragmentation functions. The newly installed Lambda Wheels, which increase the acceptance for Λ particles by about a factor of four, will allow a precise investigation of Λ electro-production and Λ polarisation effects in the current and the target fragmentation regions. A Recoil Detector, which is presently being constructed and will cover a large fraction of the solid angle around the HERMES target, will help to improve the identification of exclusive events in deeply virtual Compton scattering and hard electro-production of scalar and vector mesons. The expected data will be a first step in exploring the framework of generalised parton distributions, which is expected to expand our knowledge of hadron structure in a fundamental way as it gives access to the role of partonic correlations not contained in usual parton distribution functions.

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