Astrophysical Plasmas

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1 Introduction

99% of the luminous matter in the Universe is in plasma state, i.e. it consists of ions and electrons with some degree of ionization larger than 10^{-6}. This simple statement already proves the importance of plasma physics for astronomy. Astropyhsical plasma systems appear everywhere in space: There are extragalactic jets with spatial extensions up to some 10^5 light years. They originate from the centres of quasars, often flow with almost the velocity of light and consist of relativistic particles and pinch like helical magnetic fields. There are disk galaxies with magnetized spirals which rotate with a few hundred kms^{-1}. Supernovae explode with velocities of a few 10^4kms^{-1}, leaving a tiny nucleus of 10 km in size and of about 1.5 solar masses in form of a strongly magnetized rapidly rotating neutron star. The interstellar medium is a highly filamented magnetized medium in which shear motions compress, stretch and twist the magnetic field within a highly conducting collisionless plasma. The intergalactic medium is a very thin (1-100 particle per m^3) and very hot ($T \simeq 10^{7-8} K$) plasma. The sun is surrounded by a corona which is hundred times hotter than its surface! The Earth ionosphere glows from time to time since plasma curtains display the polar light. We really live in a plasma universe!

2 Typical parameters of astrophysical plasmas

Typical examples for astrophysical plasmas are (Peratt 1992; Tajima and Shibata 1997)
-all the stars (about 10^{11} stars per galaxy)
-the interstellar medium in galaxies (about 10^{11} objects)
-the intergalactic medium between galaxies and galaxy clusters.
The particle number density covers an enormous range from about one particle per m^3 in the intergalactic medium up to 10^{30}m^{-3} in the central regions of hydrogen burning solar like stars.
The plasma temperature ranges from only a few hundred Kelvin in cold but still significantly ionized gas clouds in which new stars form, up to a few 10^8 K in the X-ray halos of galaxy clusters that reveal themselves by thermal bremsstrahlung.
Every plasma in the universe is magnetized (e.g. Kronberg 1994)! The smallest magnetic field strengths of about 10^{-11} Tesla have been deduced from the

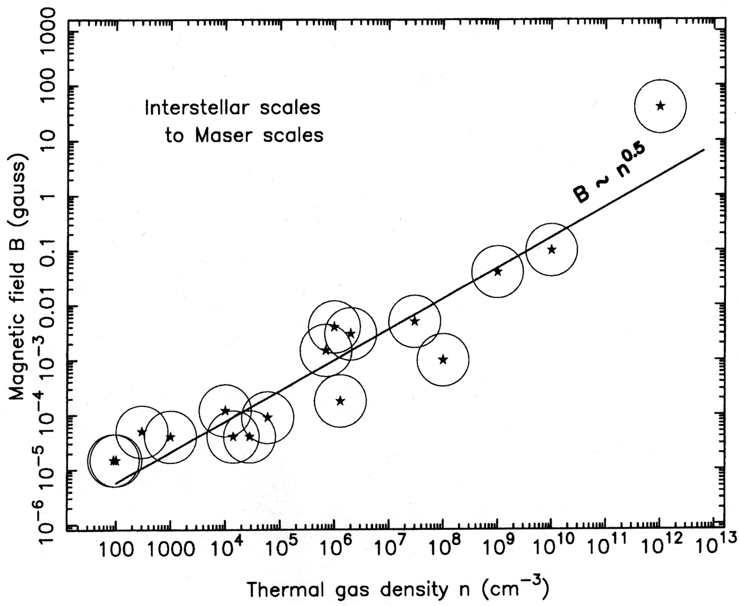

Figure 1: Observed behavior of the magnetic field B (in Gauss) as a function of the total gas density n (in cm^{-3}), for n>100 cm^{-3}. (Vallée 1997)

low- frequency radio emission in the intergalactic medium, which is due to synchrotron emission of relativistic electrons (Kim et al. 1990). The strongest magnetic fields at the surface of neutron stars are about 10^{8-10} Tesla, measured by the cyclotron emission line in the X-ray regime (Trümper et al. 1978). Neutron stars are only 10 km in size, the intergalactic medium has a typical size of millions of lightyears. The intergalactic magnetic fields are quite well ordered on spatial scales of some thousand lightyears (Kronberg 1994).
Besides the population of thermal particles, astrophysical plasmas contain ultra-relativistic particles, electrons and nuclei, as well. This population is known as cosmic rays, although it consists of particles. The relativistic electrons radiate via synchrotron emission and if they are in compact systems with intense photon fields via inverse Compton scattering (Longair 1981). We will discuss the different radiation processes in the next section. From observations of the electromagnetic spectrum we can estimate the electron energies required to provide the observed radiation. Especially from the observations from the most recent X-ray satellites CHANDRA and XMM it became clear that in AGN, in extragalactic jets (Birk and Lesch 2000)) and in supernova remnants the electrons have energies up to 100 TeV (Koyama et al. 1995; Willingdale et al. 2001) Typically at least a few hundred MeV electron population fills the space between the stars and galaxies. In most galaxies the electron population is energized globally up to 10 GeV (Lesch and Chiba 1997).
Concerning the hadronic population our knowledge is not as complete, since

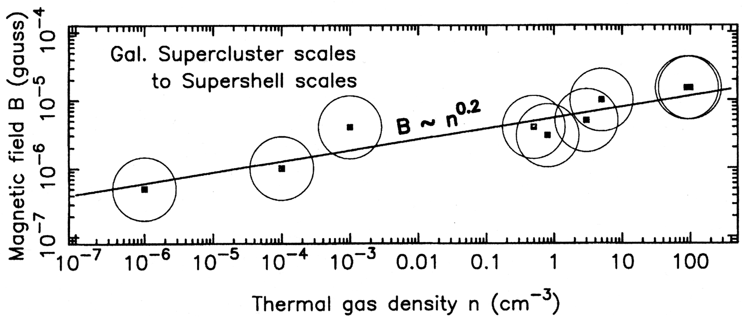

Figure 2: Observed behavior of the magnetic field B (in Gauss) as a function of the total gas density n (in cm^{-3}), for n<100 cm^{-3}. (Vallée 1997)

these particles (protons and nuclei) do not undergo intense energy losses by radiation processes like the electrons do. We know only the energy distribution of the cosmic ray hadronic population which encounters the earth. From cascades induced by cosmic rays hitting the earth atmosphere we can deduce maximum energies of about 10^{20} eV for the protons. Typical energies are in the range of MeV to GeV (Hillas 1984).
Astrophysical plasmas exhibit small scale and large scale motions at all velocities (Tajima and Shibata 1997). The plasma velocities are mainly determined by gravity and rotation. The interplay of these two forces explains the main part of plasma flows in astrophysical objects. This is true for accretion disks around black holes and neutron stars, which accelerate plasmas almost to the speed of light. Close to galactic nuclei plasma clouds move with a few 10^{3-4}kms^{-1}. In the interstellar medium clouds move with a few $10 - 100$kms^{-1}. Due to stellar explosions there are also plasma flows escaping with 20.000kms^{-1}. In any case the velocity fields are characterized by significant spatial gradients, i.e. they represent always shear motions.

After this brief introduction we can already formulate the central issues of modern plasma astrophysics:
Where do the magnetic fields come from?
What is the origin of magnetized relativistic jets?
How are the particles accelerated to the ultrahigh energies?
What is the interplay of acceleration processes and the radiative loss mechanisms?
How do the astrophysical plasmas respond to the distortions by unsaturated externally driven shear flows?
How do we have to describe such thin and highly conducting magnetized plasmas?

In the following we give an overview of the basic theoretical procedures and

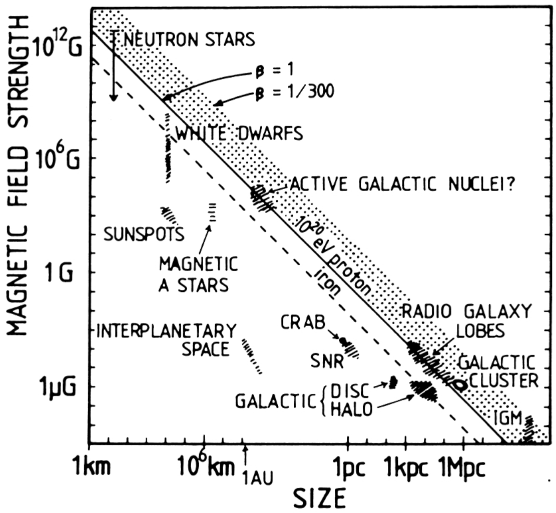

Figure 3: Size and magnetic field strength of possible sites of particle acceleration. Objects below the diagonal line cannot accelerate protons to 10^{20}eV. (Hillas 1984)

constituents necessary to tackle plasma astrophysical problems. Let us start with the source of information about the physical state of astrophysical plasmas: the electromagnetic radiation.

3 Radiation processes

(References: Rybicky and Lightman 1979, Longair 1981)
The electromagnetic radiation is the cosmic newspaper, thus, we must know what kind of mechanisms can be responsible for the photons which hit our detectors or telescopes. We observe the Universe in practically all bands of the electromagnetic spectrum starting with the radio regime with which important information can be gained about the magnetic field strength and structure since most of the radio emission in space is of nonthermal origin, i.e. it is synchrotron radiation. Thus, radio waves tell us about relativistic electrons gyrating in magnetic fields.
The infrared photons mostly originating in star-forming regions are produced by scattered light from dust particles and by direct emission due to the rotational and vibrational transitions of molecules. IR-radiation is the indicator for star formation.
Optical photons are due to stellar energy production. Optical observations performed by the largest ground based telescopes with mirrors up to 15 m and

also in space with the Hubble Space Telescope are still the backbone of modern astrophysics.
UV-photons can be observed only by satellites. They come from very hot stars and indicate young stars, since massive, hot stars which are 10-50 times heavier than the sun have lifetimes of only a few million years, whereas our sun lives 10 billion years.
The X-ray astronomy developed very rapidly in the last decade and gave deep insights into the hot plasmas in the intergalactic and interstellar medium, as well as about stellar X-ray sources and nonthermal X-ray emission from extragalactic jets.
The emission from particles with energies from the MeV- range up to 100 TeV are now also detectable (Gaidos et al. 1996). Here we gain information how particles are accelerated to the highest energies in the Universe. For example, the supernova remnant in the Crab nebula, the remnant of a supernova explosion observed by chinese astronomers in 1054, emits gamma rays at energies of about 50 TeV which demands for accelerated electrons at least at this energy level (Willingdale et al. 2001).
It is the central subject of astrophysics to translate the received electromagnetic radiation into physical models by using the theory of electromagnetic radiation. There are in principle two classes of radiation mechanisms: the thermal and nonthermal processes. The latter are mechanisms like synchrotron radiation and inverse Compton scattering, where relativistic electrons emit photons due to the interaction with magnetic fields or external radiation fields. Nonthermal radiation is produced by particles with nonthermal energy distributions, mainly with power law distribution functions or almost monoenergetic distributions (Lesch 2000).
In the case of thermal radiation the energy distribution of the particles is a Maxwellian which defines a temperature. The particles will either emit lines (atomic transitions) or continuum radiation (free-free emission, i.e. bremsstrahlung). From the blue and/or redshifts of the recombination lines we get knowledge about the plasma velocities and densities. We will not discuss the numberless line emission or absorption mechanisms but rather concentrate on the emission of a fully ionized plasma without recombination lines, etc., and we will also not consider black-body radiation, which is well-known (Rybicki and Lightman 1979 for a review of radiation transport) This leads us to the mechanism of thermal bremsstrahlung which gives information about the plasma temperature T and the electron density n_e.

3.1 Thermal Bremsstrahlung

When a charged particle is accelerated or decelerated it emits electromagnetic radiation. Bremsstrahlung is the kind of radiation emitted in all electromagnetic encounters between the charge and the nuclei of the medium through which it passes. Of special importance is the thermal bremsstrahlung, where free electrons with a kinetic energy $m_e v_e^2/2$ equal to the thermal energy $k_B T_e$ (k_B denotes the Boltzmann constant) move in a hydrogen plasma. The radiated energy is a direct measure of the plasma temperature and the plasma density, since the energy loss rate of the plasma is

$$\left(\frac{dE}{dt}\right)_{brems} = 1.435 \times 10^{-38} T_e^{1/2} n_e n_i \;\; [Wm^{-3}] \qquad (1),$$

with the electron temperatur T_e in eV. For quasineutral plasmas with single charged ions, i.e. $n_e = n_i$, bremsstrahlung measures the electron temperature and electron (ion) density.

The bremsstrahlung spectrum is characterized by a flat curve in the optically thin regime followed by an exponential break at $h\nu \simeq k_B T_e$ (where h denotes the Planck constant). In the optically thick regime the spectrum rises proportional to ν^2.

Thermal bremsstrahlung is an important radiation process for very hot gases radiating thermal X-rays (like supernova remnants and gas halos of galaxy clusters). Optically thick bremsstrahlung with a flat spectrum is emitted in the radio regime by HII-regions that are located close to the regions of star formation.

3.2 Synchrotron Radiation

Synchrotron emission is the radiation of a high energy particle gyrating in a magnetic field. It was originally noted in some early betatrons where high energy particles were first accelerated to ultrarelativistic energies. This same mechanism is responsible for the radio emission from the Galaxy, from supernova remnants and extragalactic radio sources. It is also responsible for the nonthermal optical and X-ray emission in the Crab Nebula and for the optical continuum emission of quasars and jets. The sun and young stellar objects show radio flares.

For nonrelativistic velocities the complete nature of the radiation is simple and is called cyclotron radiation. The frequency of the emission is the frequency of the gyration of charged particles in magnetic fields

$$\omega_g = \frac{zeB}{m_e} \qquad (2)$$

where B and z denote the magnetic field strength and the charge number. A useful figure to remember is the gyrofrequency of an electron

$$\nu_g = \frac{eB}{2\pi m_e} = 28\,\mathrm{GHz}\; B[T]. \qquad (3)$$

Particles radiate only, if their pitch angle θ (which is the angle between the directions of the particle velocity and the magnetic field line) is not zero. The energy loss rate of an electron

$$-\left(\frac{dE}{dt}\right) = 2\sigma_T c \gamma^2 \frac{B^2}{2\mu_0} sin^2\theta \qquad (4)$$

where $\sigma_T = 6.65 \times 10^{-29} m^2$ denotes the Thomson cross section and γ is the Lorentz factor of the particles. This result applies for electrons of a specific pitch angle θ. To get the average loss rate for particles of all pitch angles, we must average over solid angle, i.e. $P(\theta)d\theta = \frac{1}{2} sin\theta d\theta$,

$$-\left(\frac{dE}{dt}\right)_{average} = 2\sigma_T c\gamma^2 \frac{B^2}{2\mu_0}\frac{1}{2}\int_0^\pi sin^3\theta d\theta = \frac{4}{3}\sigma_T c\gamma^2 \frac{B^2}{2\mu_0}. \tag{5}$$

The typical frequency emitted by an electron is the nonrelativistic gyrofrequency times γ^2

$$\nu_{syn} \simeq \gamma^2 \nu_g sin\theta. \tag{6}$$

The emitted radiation spectra depend on the energy distribution of the relativistic electrons. An almost monoenergetic (beam- like) distribution of relativistic electrons emit an optically thin synchrotron spectrum

$$I(\nu) \propto \nu^{1/3} \;\; (\nu \ll \nu_{syn}) \tag{7}$$

and

$$I(\nu) \propto exp\left(-\frac{\nu}{\nu_{syn}}\right) \quad (\nu \gg \nu_{syn}) \tag{8}$$

with a maximum frequency $\nu_{max} \simeq 0.29\nu_{syn}$. Such a spectrum has been detected in several galactic nuclei and especially in the centre of our galaxy (Sgr A*) and its immediate surrounding – the arc. There the synchrotron emission is concentrated in long (150 light years) and thin (0.5 light years) highly polarized magnetic filaments. The magnetic field strengths in these filaments is estimated to be about 10^{-7} T. Given a measured maximum frequency of about 100 GHz this indicates that some process has produced an electron energy distribution which is almost mononenergetic with a characteristic energy of about 40 GeV (Lesch and Reich 1992). In the central light year of the galactic centre, in SgrA*, the monoenergetic electrons have energies of some MeV (Lesch et al. 1988, Duschl and Lesch 1994).
Nevertheless, such monoenergetic spectra are rare. More typical are energy distributions of power law form

$$N(E)dE = \kappa E^{-x} dE \tag{9}$$

where$N(E)dE$ refers to the number of particles per unit volume. Such energy spectra emit optically thin radiation power laws

$$I(\nu) \propto \nu^{-(x-1)/2} \tag{10}.$$

The important rule is that *if the electron energy spectrum has a power law index x, the spectral index of the synchrotron emission of these electrons is* $\alpha = (x-1)/2$. Such spectra are found in supernova remnants with $\alpha \simeq 0.5$

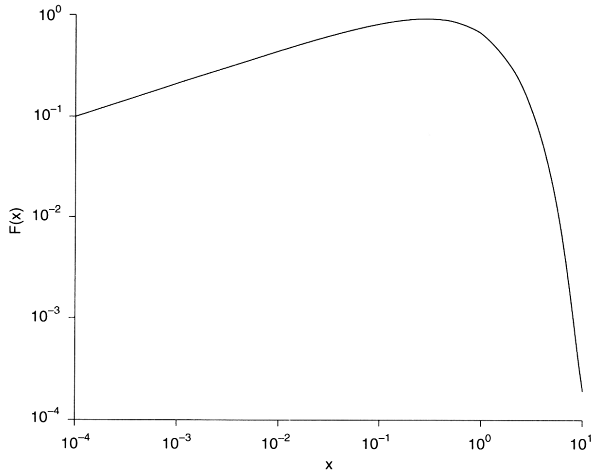

Figure 4: The intensity spectrum of the synchrotron radiation of a single electron with logarithmic axes. The function is plotted in terms of $x = \nu/\nu_{syn}$.

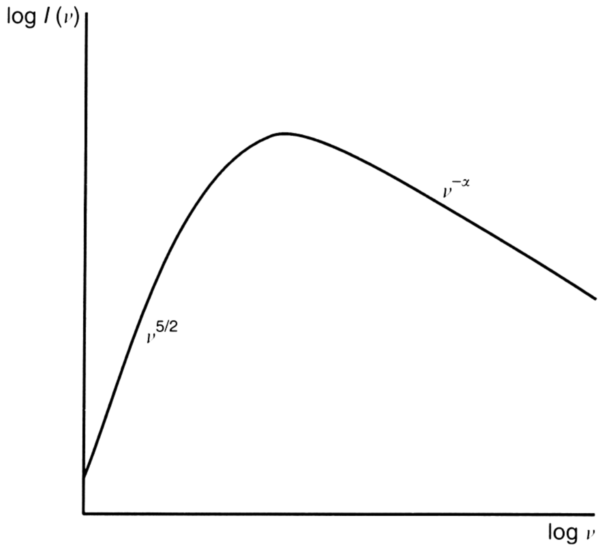

Figure 5: The spectrum of a source of synchrotron radiation with a power-law energy distribution of the relativistic electrons. At low frequencies the phenomenon of self-absorption is exhibited.

indicating an energy distribution of E^{-2}, in galaxies with $\alpha \simeq 0.83$, i.e. $E^{-2.66}$ and in galactic nuclei with rather flat spectra of about E^{-1} (Birk et al. 2001).

Synchrotron radiation is linearly polarized perpendicular to the magnetic field. Its degree of polarization in case of a power law is

$$\Pi = \frac{x+1}{x+\frac{7}{3}} \tag{11}$$

giving a maximum linear polarisation of about 72%.

The optically thick synchrotron spectrum does not depend on the energy distribution and has the form

$$I(\nu) \propto \nu^{5/2}. \tag{12}$$

Such optically thick synchrotron spectra are observed in nuclei of quasars which evolve to a $\nu^{-\alpha}$ optically thin spectrum at higher frequencies.

3.3 Inverse Compton Scattering

Relativistic electrons loose energy when they collide with photons, because the photons are scattered to much higher energies. This process is called *inverse Compton scattering*, because in the usual Compton scattering interaction the photon gives energy to a particle initially at rest and thereby looses energy in the collision. The situation most often encountered in astrophysical plasmas is the limit in which the energy of the photon in the centre of momentum frame of the collision is much less than the electron's rest mass energy, i.e.$\gamma h\nu \ll m_e c^2$. The energy loss rate due to inverse Compton scattering is

$$\left(\frac{dE}{dt}\right)_{IC} = \frac{4}{3}\sigma_T c \gamma^2 c U_{rad} \tag{13}$$

where U_{rad} denotes the energy density of the external photon field. Note the great similarity of this result to that for synchrotron radiation all the way down to the factor of 4/3. The reason for this is that in both cases the particles are accelerated by the electric field no matter where the field originates from. Due to this similarity it is clear what the typical emitted frequency in inverse Compton scattering must be

$$\nu_{IC} \simeq \gamma^2 \nu_0 \tag{14}$$

where ν_0 denotes the photon frequency of the external radiation.
In the limit $\gamma h\nu \ll m_e c^2$ inverse Compton scattering is Thomson scattering. If instead the Doppler shifted photon energy $\gamma h\nu$ in the rest frame of the electron is equal to the electrons rest mass energy $m_e c^2$ the cross section is given by the Klein Nishina cross section which is approximately $\sigma_{KN} \propto \sigma_T/\gamma$.
The IC-process is important in compact objects like AGN, in which very intense radiation fields interact via Compton scattering with accelerated plasma flows in accretion disks around black holes with masses of a few million solar masses. The

disks are heated to temperatures of about 10^5 K, which can be deduced from line emission and thermal continuum radiation. Such disks emit intense radiation in the optical and UV at about 10^{13-15} Hz. Some AGN are intense emitters of TeV-radiation, i.e. 1000 GeV! The only mechanism to produce such a high energy radiation without intense pair production (remember the rest mass of an electron plus a positron is about 1MeV) is optically thin inverse Compton scattering. One requires electrons with energies of several TeV, their corresponding Lorentz factors are about 10^6 which adds up to $\nu_{IC} \simeq \gamma^2 \nu_0 \simeq 10^{25}$Hz! This scattering works efficiently only for the low energy photons of the external radiation. For UV photons the Klein-Nishina limit holds and the cross section drops with $1/10^6$, since the photon energy in the rest frame of the electrons is comparable to the rest mass energy of the electrons.
Finally we give some interesting figures for the loss time scales of inverse Compton scattering and synchrotron radiation For the IC-case we have a loss time

$$t_{loss}^{IC} = \frac{5 \times 10^{21}\mathrm{sec}}{\gamma U_{rad}} \tag{15}$$

with U_{rad} in J/m^3.

For synchrotron radiation we have

$$t_{loss}^{SYN} = \frac{3 \times 10^{7}\mathrm{sec}}{\gamma B^2} \qquad (16),$$

with B in 10^{-4}T.

For more information about radiation processes we refer to Rybicki and Lightman (1979).

4 Cosmical Magnetohydrodynamics

(Reference: Tsinganos 1996)
Since astrophysical plasmas are highly conductive, the ideal magnetohydrodynamics (MHD) may be an appealing description for such magnetized gases. For large spatial scales this statement is correct, i.e. in cosmical plasmas magnetic fields obey the frozen in condition which means that the magnetic flux $\Psi = BR^2$ is conserved globally. In terms of Ohm's law an ideal MHD-plasma obeys $\mathbf{E} = -\mathbf{v} \times \mathbf{B}$, i.e. induced electric fields force plasma and magnetic field to corotation and frozen-in flux. Large volumes in space like the interstellar medium in galaxies, the intergalactic medium or extragalactic jets can be almost perfectly described by ideal MHD.

The ideal description, however, is not valid locally, i.e. for small spatial scales, in which magnetic boundary layers are build up by external forces like gravity, rotation or explosions. There, the magnetic field is sheared, twisted and compressed leading to locally antiparallel field structures, that is to the formation of current sheets. In such transition regions the ideal nondissipative MHD-description cannot be applied to cosmical plasmas. In the boundary layers often such strong electrical currents flow that collective processes provide a local reduction of the conductivity and that rather fast energy dissipation sets in via magnetic reconnection. According to Ohm's law the deviation from idealness is equivalent to a parallel electric field which may accelerate plasma particles to ultrarelativistic energies as it is required by the observations.

Let us start with the ideal MHD description and its consequences for the dynamics of astrophysical plasmas.

4.1 The Ideal MHD Equations

One important equation of ideal MHD describes the time evolution of magnetic fields due to induction ($\mathbf{v} \times \mathbf{B}$). It can be derived from the Maxwell equation $\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t}$ and Ohms law $\mathbf{E} = -\mathbf{v} \times \mathbf{B}$, leading to

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}). \tag{17}$$

According to this equation it is only the induction which changes the field strengths with time. This equation is synonymous with the frozen-in condition of the magnetic flux.

An equation which has profound implications for the energy transport in ideal MHD systems like astrophysical plasmas is

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{j} \tag{18}$$

which can be rewritten as

$$\nabla \cdot \mathbf{j} = 0. \tag{19}$$

In ideal MHD-plasmas the displacement-term $\frac{1}{c^2}\frac{\partial \mathbf{E}}{\partial t}$ can be neglected, because the conductivity is assumed to be so high that electric fields are always much weaker than magnetic fields, i.e. $E \ll B$, and all velocities are smaller than the velocity of light. The condition that the divergence of the current density vanishes means that the current circuit has to be closed. If we distinguish between currents running parallel ($\mathbf{j}_{||}$) and perpendicular ($\mathbf{j}_{\perp}$)to the magnetic field Eq. (19) represents an important relation:

$$\nabla \cdot \mathbf{j}_{||} = -\nabla \cdot \mathbf{j}_{\perp}. \tag{20}$$

Eq. (20) means that the perpendicular currents are sources for parallel currents and vice versa! If, for example, in a magnetized plasma forces are present which drive plasma motions perpendicular to the magnetic field electromagnetic energy is stored in the plasma. The region where this happens can be regarded as MHD generator. It is connected with the outside plasma by magnetic field-aligned currents, which are force-free (supposed that the external medium is of lower density, i.e. the pressure forces outside of the generator are weaker than inside the generator). Force-free currents are perfect for energy transport! They usually result in a filamentary magnetic field structure which stabilizes the plasma flow along the magnetic field and allows energy deposition far away from the generator. Such a scenario depends on the pressure gradients which may stabilize a given magnetic configuration. In the denser generator region pressure gradients will be definitely important.

In the external medium, which may be the corona of a star, the halo of a galaxy or an accretion disk around a black hole, the pressure is weak and the pressure gradients are small. The strong shear of the magnetic fields in the generator may be transported into that external medium by Alfven waves, i.e. the magnetic field strengths can be enhanced by the kinetic energy of the plasma in the generator. The weakness of the pressure and the stronger magnetic field both support the development force-free filaments which are energetically preferred stages of magnetized plasmas in which currents are driven. Such force-free filaments are characteristic features of any astrophysical plasma with low pressure (Tajima and Shibata 1997)

In any case at least partially the stored energy in the filaments will be dissipated by either plasma heating or particle acceleration. Both mechanisms should be observable in terms of thermal radiation in heated plasmas (like the quiet solar corona) or in terms of nonthermal radiation where particles have been accelerated (like in solar flares). However, the dissipation of the energy is not described by ideal MHD. We will discuss that important subject in the subsequent sections.

Here we list typical examples of configurations in which an external force agitates a plasma perpendicular to the magnetic field lines, stores energy and transports this stored energy via magnetic field-aligned electric currents, thereby inducing magnetic filaments:

Solar and young stellar flares - solar prominences - cicumstellar disk - Earth and Jovian aurora - interstellar filaments - stellar jets - extragalactic jets - accretion disks around AGN.

This list is time dependent in that sense that it depends on the actual state of the observational techniques in astrophysics. The higher the spatial resolution of the instruments is, the more spatial structures, i.e. inhomogenities, are detected. The more advanced astrophysical instruments are, the more plasma physics becomes relevant for the interpretation of the observations!

For astrophysicists the plasma structure and the magnetic field structure is of course important for the understanding of stability and inhomogeneity. However, as mentioned above, the most important subject of plasma astrophysics concerns the electromagnetic radiation which can be expected from dynamical plasma processes in terms of plasma heating and particle acceleration. To investigate such problems we have to extend ideal MHD to resistive MHD that includes finite electrical conductivity and thereby energy dissipation.

4.2 The resistive MHD Equations

The MHD equations that take into account the role of a finite electrical conductivity σ are quite different from the set of ideal MHD equations in many ways. Most important, Ohm's law now reads

$$\mathbf{j} = \sigma(\mathbf{E} + \mathbf{v} \times \mathbf{B}). \tag{21}$$

The electrical conductivity $\sigma isgiven$ in its most simplest form by

$$\sigma = \frac{n_e e^2 \tau_{ei}}{m_e} \tag{22}$$

where $\tau_{ei} \simeq 0.266 \times 10^6 T_e^{3/2}/n_e \ln\Lambda$ denotes the collision time (with the Coulomb logarithm $\ln\Lambda$) which is generally between 5 and 20 and has a weak dependence on temperature and density). In case of Coulomb collisions the conductivity is $\sigma = 1.53 \times 10^{-2} T^{3/2}/n_e \ln\Lambda$. For temperatures higher than 10^3 K the conductivity is so high that the frozen-in condition is fulfilled. Nevertheless even in case of only Coulomb collisions the conductivity is finite and has to be taken into account if the plasma is known to be dissipative. Finite conductivity effects become important for the plasma dynamics in boundary layers. This is comparable to the influence of finite viscosity in thin sheets around obstacles in almost ideal hydrodynamic flows. The ideal induction equation (17) changes now into

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}) - \nabla \times (\eta \nabla \times \mathbf{B}) \tag{23}$$

where $\eta = 1/(\mu_0 \sigma)$ is the *magnetic diffusivity* which must not be confused with the electrical resistivity which is just the inverse of the electrical condcutivity. If the magnetic diffusivity is uniform the we get the equation usually used as induction equation

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}) + \eta \nabla^2 \mathbf{B}) \qquad (24).$$

Eq.(24) for a vanishing plasma velocity **v** describes the pure diffusion of magnetic field lines due to finite resistivity/conductivity. Eq.(23) with a spatial dependent diffusivity describes localized dissipation of magnetic field energy. Since most astrophysical plasmas are globally ideal and dissipative only on small spatial scales, we will continue with the consequences of Eq. (23). What does this equation mean? According to the non-ideal Ohm's law a localized resistivity corresponds to a localized electrical current distribution. Strong currents imply strong gradients in the magnetic fields that are induced by the currents. As long as only Coulomb collisions are the source for resistivity the field lines almost perfectly co-move with the plasma, since collisions are very rare in typical astropysical plasma systems. However, if the magnetic field lines are subject to strong plasma motions, the situation changes drastically. The frozen-in field lines may be twisted, stretched, compressed etc.., and thereby strong currents are induced. Especially in the case of shear motions (like in the case of differential rotation), or turbulent gas flows, the magnetic energy can very efficiently be dissipated via **magnetic reconnection** in very localized current sheets.

Reconnection is an intrinsic property of a magnetized plasma with shear and/or turbulent motions. The encounter of magnetic field components with different polarity corresponds to parallel electric currents which attract each other. In that sense any kind of plasma motion can trigger reconnection, provided that any localized violation of ideal Ohm's law allows for changes of the magnetic field topology.

If the plasma pressure between opposite fields $\pm B$ is insufficient to keep the fields apart (e.g. by pushing the different flux systems apart), the plasma squeezed from between them and the two fields approach each other. The field gradient steepens and eventually the current density $\sim \nabla \times B$ becomes so large that there is strong current driven dissipation. One crucial problem in the context of magnetic reconnection concerns the nature of the dissipation of the currents with density $j = en_{\rm e}v_{\rm d}$ is provided (e is the charge and $v_{\rm D}$ denotes the drift velocity of the electrons). Strong dissipation does not exclusively imply an enhancement of the current density but also can be caused by a localized reduction of the conductivity. Indeed both effects support each other. This can be understood as follows: when the conductivity is high, the magnetic field lines are "frozen into" the plasma. Field lines embedded within a volume element of plasma are carried along by the moving plasma. Any two plasma elements that are threaded by the same lines of force will remain threaded in this way, i.e. two plasma elements that are threaded by different lines of force can never be magnetically linked together. Thus a high conductivity prevents the field lines from merging. In astrophysical plasmas reconnection occurs only in those plasma regions where the electrical conductivity is drastically reduced below its classical Coulomb value. The reduction of the electrical conductivity can be provided by plasma microinstabilities that are driven by the induced currents. If the plasma is locally unstable, i.e. if the current density exceeds a critical value $j > j_{\rm crit}$ or $v_{\rm d} > v_{\rm crit}$ (e.g. Huba 1985) microinstabilities will be excited. In the non-linear saturation states microturbulent electromagnetic fields enhance the collision frequency by wave-particle interactions; they lead via *anomalous conductivity* by momentum tranfer between the charged particles via the microturbulent electromagnetic fields. The effective σ locally becomes anomalously low, greatly enhancing the dissipation and reconnection of the lines of force. Then the concept of "frozen in" field lines is no longer valid in the whole plasma and the plasma locally moves relative to the field lines. When this occurs, strong electric fields along the reconnection length L are induced. By such electric fields the magnetic energy is converted to particle energy.

Schindler et al. (1988, 1991) have shown that in three dimensions the reconnection process is always related to a magnetic field aligned electric potential $U = -\int E_{\|} ds$. The induced electric field is always parallel to that magnetic field component which is not directly involved in the reconnection process, i.e. that is directed perpendicular to the plane of antiparallel magnetic field components. This result was used to consider the acceleration of particles in cosmical plasmas (e.g. Birk et al. 2001) and it was concluded that parallel electric fields associated with nonideal plasma flows can play an important role in cosmic particle acceleration (Schopper et al. 2002). The actual value of $E_{\|}$ depends on the details of the microscopic instability, which is responsible for the deviations of the plasma from the ideal high conductivity state. It has been shown that in the case of the slow reconnection mode first proposed by Parker (Parker 1979) it is possible to get an estimate of the maximum Lorentz factor particles can attain via parallel electric potentials (e.g. Lesch and Reich 1992).

Starting from the assumption of a stationary Ohmic dissipation in a two-dimensional reconnection sheet with an area $\sim L^2$ and a thickness l, the dissipation surface density in the sheet lj^2/σ is just to devour the influx of magnetic energy $uB_0^2/2\mu_0$ where u is the velocity of the approaching field lines. Conservation of fluid mass density requires that the net magnetic field inflow balances the outflow

$$uL = v_A l. \tag{25}$$

In terms of the magnetic Reynolds number $R_{\mathrm{M}} = 2Lv_{\mathrm{A}}/\eta$ one gets (Parker 1979)

$$l = \frac{2L}{\sqrt{R_{\mathrm{M}}}} \tag{26}$$

and

$$u = \frac{2v_{\mathrm{A}}}{\sqrt{R_{\mathrm{M}}}} \tag{27}$$

The thickness of a reconnection sheet is also defined by Maxwell's equation $\nabla \times \mathbf{B} = \mu_0 \mathbf{j}$, which means $l \simeq B/\mu_0 j$. Now the whole problem is shifted to the microscopic level of description. As mentioned above the particle collisions often are not efficient enough to produce a significant magnetic diffusivity (or to decrease the electric conductivity). Current driven plasma instabilities are involved. The resulting anomalous collision frequency depends on the drift velocity between electrons and ions. We choose in the following the wave mode which has the lowest instability threshold where the drift velocity v_{d} is equal to the ion thermal velocity $v_{\mathrm{thi}} = \sqrt{k_{\mathrm{B}}T_{\mathrm{i}}/m_{\mathrm{i}}}$ namley *the lower hybrid (LH) wave* with the frequency $\omega_{\mathrm{LH}} \simeq 1.26 \cdot 10^{10} B$. LH waves depend on the magnetic field strength, which means they are also suitable for three-dimensional reconnection where the magnetic field is not zero in the reconnection zone. A fully developed LH-instability results in an effective collision frequency of the order of ω_{LH} (Sotnikov et al. 1978). We note that the central role of LH-waves in reconnection zones has been established for magnetospheric activity (e.g. Shapiro et al. 1994).

The parallel electric field component cannot exceed the perpendicular one. Inserting u into the electric field $E = \mathbf{v} \times \mathbf{B} \sim \mathbf{uB}$ and l into L, using $\nu_{\mathrm{coll}} \simeq \omega_{\mathrm{LH}}$ and with

$$\gamma m_{\mathrm{e}} c^2 \simeq eEL \tag{28}$$

one obtains the maximum Lorentz factor electrons can achieve in a three-dimensional reconnection zone, where the conductivity is determined via lower-hybrid waves (Lesch 1991)

$$\gamma_{\mathrm{LH}} \simeq 6 \cdot 10^7 B \simeq 6 \cdot 10^5 \left[\frac{B}{0.01\mathrm{T}}\right]. \tag{29}$$

In the last years we have performed MHD numerical simulations to study the reconnection dynamics (Lesch and Birk 1997, 1998; Birk 1998; Birk et al. 1998; Konz et al. 2000). Using a self-consistent description of the localized reduced current dependent conductivity and fully relativistic test-particle simulations the effiency of particle acceleration to relativistic energies, including energy losses via synchrotron and inverse Compton radiation, has been proven (Schopper et al. 1998, 1999).
Our results can be summarized within one sentence: *Magnetic reconnection is a very efficient acceleration mechanism.* In reconnection regions due to sheared magnetic fields the radiation losses in many astrophysical contexts do not influence the acceleration, simply because the acceleration is much faster than the losses. Depending on the initial injection position particles experience different electric field strength and power laws appear. Depending on the isotropy of the injection spectrum, a significant fraction of the low-energy particles which enter the reconnection region are accelerated. In case of an anisotropic injection, the resulting energy distribution function of the electrons will be quasi-monoenergetic, since the achievable energy depends on the length of the reconnection zone as in a linear accelerator. Since the final energy distribution will show an accumulation of particles either at the maximum energy or at the energy where radiation losses and acceleration exactly cancel, we obtain an energy distribution which exhibits pronounced low and high energy cutoffs, i.e. a relativistic electron beam (REB) (Lesch and Schlickeiser 1987; Birk et al. 2001).

In the context of flares from AGN, the energy given by Eq. (29) is close to the required energies for TeV- radiation (Birk et al. 1999). The particles would not lose their energies completely via synchrotron radiation in the strong magnetic field, because they flow along the field lines, i.e. the energy distribution is highly anisotropic (Crusius-Wätzel and Lesch 1998).

In this section we have used conventional reconnection in terms of an enhanced collision frequency via plasma waves, excited by some phase-space instability. In a rarefied and relativistic plasma, as can be found in extragalactic jets, the excitation of waves is much more difficult. The plasma is very stable due to the high energy of the particles. That is, if we still have external shear flows which twist and compress the magnetic field lines, thereby enhancing locally the current density the system starts to become *filamentary* (Wiechen et al. 1998; Lesch and Birk 1998, Konz et al. 2000, Birk and Lesch 2000). The plasma is now completely collisionless, there are even no plasma fluctuations to reduce the conductivity, i.e. we have no dissipation at all. In such circumstances, like for example in jets, where the plasma density is very low and only magnetic fields and relativistic particles react to the externally driven differential rotation of the disk, the system reacts via strong filamentation and the inertia of the particles play a very important role. We describe *the inertia- driven -reconnection* in the next section.

5 Inertia-driven reconnection in AGN

How reacts an ideal collisionless, magnetized plasma which is strongly distorted by externally driven shear flows? The magnetic fields will be twisted until tube like structures appear. On the same spatial scales the magnetic field gradients will become very strong, i.e. current sheets form and magnetic reconnection sets in. The magnetic field aligned electric field that appears in a reconnection zone is driven by the reduction of the local electrical conductivity. In principle, the strength of such a field cannot exceed the convective electric field, which is responsible for the shear, twist, etc.., given by Eq.(28). The maximum shear velocity is given by the Alfvenic speed $v_A = B/\sqrt{\mu_0 \rho}$.
Thus, we get

$$E_{\|} = \zeta v_A B \tag{30}$$

where ζ denotes the effiency of the conversion of the generated electric field energy to particle energy.
The necessary condition for magnetic reconnection to operate is the local violation of the ideal Ohm's law

$$\mathbf{E} + \mathbf{v} \times \mathbf{B} = \mathbf{R}, \tag{31}$$

where $\mathbf{R} \neq 0$ is some unspecified nonideal term that gives rise to a nonvanishing electric field component parallel to the magnetic field.
Since the relativistic plasmas in AGN and especially inside the jets from AGN are highly collisionless, i.e. the dissipative term in Equation (21) is negligible. Therefore, the question arises concerning the nature of the nonidealness and also about the effectiveness ζ in a completely collisionless plasma. Rarefied high energy plasmas are very stable against the excitation of plasma waves, since the wave excitation condition $v_{drift} > v_{the} \sim c$ is hardly to fulfill for relativistic energies.
If a plasma is completely collisionless only the particle inertia is left to introduce some resistance. With a time dependent current density Ohm's law reads as

$$\frac{m_e}{n_e e^2} \frac{\partial \mathbf{j}}{\partial t} = \mathbf{E} - \frac{\mathbf{j}}{\sigma_{coll}}, \tag{32}$$

where σ_{coll} denotes an assumed collisional condctivity.

From dimensional analysis we see from Equation (32) that the inertial term is of order

$$\frac{m_e}{n_e e^2} \frac{\partial \mathbf{j}}{\partial t} \simeq \frac{m_e}{n_e e^2} \frac{\mathbf{j}}{\tau}, \tag{33}$$

where τ is the characteristic time of the system. Thus the inertial term can be large compared to the collisional term if $\tau \ll \tau_{coll}$, where τ_{coll} is the collisional time scale. We can solve Equation (33) explicitely (assuming zero initial current)

$$j = \sigma_{coll} \left[1 - exp\left[\frac{-\tau}{\tau_{coll}} \right] \right] E. \tag{34}$$

We therefore define an effective conductivity due to the inertia of the particles σ_i, as

$$\sigma_i = \sigma_{coll} \left[1 - exp \left[\frac{-\tau}{\tau_{coll}} \right] \right] . \tag{35}$$

The time τ in Equation (35) is regarded as the lifetime of the particle in the system, that is the time during which the electric field can be effective in accelerating the particles. If $\tau \gg \tau_{coll}$ the inertial conductivity σ_i approaches the collisional conductivity σ_{coll}. In a system with $\tau \ll \tau_{coll}$ in which the characteristic time scales are much smaller than the collision time, i.e. in thin relativistic plasmas for example, one finds

$$\sigma_i(\tau) = \frac{\sigma_{coll}\tau}{\tau_{coll}} = \frac{ne^2\tau}{m_e} \tag{36}$$

and σ_i *becomes arbitrarily small for* $\tau \to 0$. Thus, the lifetime of the particle in the system replaces τ_{coll} in the conductivity expression leading to this effective conductivity. Consequently, the inertial term in Ohm's law, usually ignored in MHD, becomes important when the lifetime of the particle in the system is small compared to a mean time between collisions.
The field aligned electric field supported by electron inertia is (e.g. Vasyliunas 1975)

$$E_{\|} \simeq \frac{m_e}{ne^2}\frac{\partial j}{\partial t} \simeq \frac{1}{\omega_{pe}^2}\frac{B}{L_{shear}\tau_A} = \frac{\lambda_{skin}^2}{L_{shear}^2} v_A B, \tag{37}$$

where L_{shear} and λ_{skin} denote the thickness of the current sheets of the filamentary current carrying magnetic flux tubes and the electron skin length, respectively. We now have a measure for the conversion effiency ζ in Equation (30)

$$\zeta \propto \frac{\lambda_{skin}^2}{L_{shear}^2}. \tag{38}$$

Since there are no collisions, the only characteristic time scales involve energy losses by radiation. Promising applications of these considerations are flares in active galactic nuclei (AGN) and extragalactic jets. For the central part we have strong inverse Compton losses (Birk et al. 1999) and for regions further away from the center and along the jets it is synchrotron radiation (Lesch and Birk 1998). The latter case allows a satisfying explanation for extended optical synchrotron emission in jets (Meisenheimer 1996) which requires continuous reacceleration along the jet flow until $\gamma \sim 10^6$!. In the case of the Centaurus A jet even synchrotron X-ray emission can result from inertia driven reconnection sites(Birk and Lesch 2000). The first case gave us preliminary hints toward a self-consistent accretion model for particles in the very centre of an active galactic nucleus towards TeV-energies.

6 Conclusions

Finally, some general comments may be appropriate. The fact that AGN are magnetized plasmas containing relativistic and nonrelativistic plasma flows which transport energy from a galactic nucleus into the jets, hot spots and lobes, request for a nondissipative energy transport. This suggests that force-free magnetic structures are very important for the plasma dynamics in AGN. Force-free means that no Lorentz force is exerted onto the plasma since the currents are "field- aligned", i.e. they flow parallel to magnetic field lines: $\mathbf{j} \times \mathbf{B} = \nabla \times \mathbf{B} \times \mathbf{B} = 0$. This condition can be satisfied in three ways: $\mathbf{B} = 0$, $\nabla \times \mathbf{B} = 0$, i.e. $\mathbf{j} = 0$ or $\nabla \times \mathbf{B} = \phi \mathbf{B}$, where the scalar $\phi = \phi(r)$ in general. Such configurations tend to have a twisted or sheared appearance and represent the lowest state of magnetic energy that a closed system can achieve (Taylor 1986). This has two important consequences. It proves the stability of force-free fields with constant ϕ and shows that in a system in which the magnetic forces are dominant and in which there is a mechanism to dissipate the fluid motion, force-free fields with constant ϕ are the natural final configuration.

The advantages of force-free magnetic field systems are obvious: Force-free structures are globally stable and thus allow an almost unperturbed transport of energy from the generator region up to some volume in which dissipation takes place (only shearing instabilities at the edge of plasma streams with the ambient medium will distort the energy transport) until the jet energy is dissipated in the hot spots. The internal energy overshoot of the plasma in such a flux tube is disspated via continuous particle acceleration. Since the currents flow parallel to magnetic field lines they will evolve into a filamentary structure simply by inducing local toroidal magnetic fields which isolate one filament from the other. In the context of acceleration it is then important to recognize the fact that the jets are collisionless. A probable shear of current carrying filaments cannot be opposed by some resistivity since the particles do not collide. The current filaments shrink to a size where the particle inertia becomes important, i.e. to filaments with thicknesses of the order of the electron skin length. It is within these filaments that the kinetic energy of the shear is dissipated via particle acceleration as described in our present paper.

To summarize, the proposed highly filamentary structure of the plasma in AGN is a natural configuration of magnetohydrodynamic flows under the influence of external forces. The lowest energy state of force-free magnetic fields allows a simple explanation for the continuous stable jet flows. Inertia-driven magnetic reconnection causes the localized dissipation of magnetic energy and in situ particle acceleration in the thin filaments all along the jet. Radiation losses of the relativistic particles only can limit the acceleration.

7 Summary

Plasmaphysics is a fundamental tool for the investigation of astronomical objects. Although gravity is the dominant force on large spatial scales the elec-

tromagnetic interaction of the ionized gases determines in numberless cases the dissipative properties which are observable as electromagnetic radiation.
The most recent progress in observational astronomy proves the necessity to consider nonlinear plasma processes like particle acceleration and plasma heating in terms of mechanisms like magnetic reconnection, in which localized regions convert magnetic energy partially into relativistic particles and/or hot plasmas.

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