58/95 JS(2)

PREPARED FOR THE U.S. DEPARTMENT OF ENERGY,
UNDER CONTRACT DE-AC02-76-CHO-3073

PPPL-3032
UC-427

PPPL-3032

IMPORTANT PLASMA PROBLEMS IN ASTROPHYSICS

BY

R.M. KULSRUD

JANUARY 1995

PPPL

PRINCETON
PLASMA PHYSICS
LABORATORY

PRINCETON UNIVERSITY, PRINCETON, NEW JERSEY

NOTICE

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial produce, process, or service by trade name, trademark, manufacturer, or otherwise, does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

NOTICE

This report has been reproduced from the best available copy.
Available in paper copy and microfiche.

Number of pages in this report: 45

DOE and DOE contractors can obtain copies of this report from:

Office of Scientific and Technical Information
P.O. Box 62
Oak Ridge, TN 37831;
(615) 576-8401.

This report is publicly available from the:

National Technical Information Service
Department of Commerce
5285 Port Royal Road
Springfield, Virginia 22161
(703) 487-4650

DISCLAIMER

Portions of this document may be illegible in electronic image products. Images are produced from the best available original document.

Important Plasma Problems in Astrophysics

Russell M. Kulsrud

Princeton Plasma Physics Laboratory and
Princeton University Observatory, Peyton Hall
P.O. Box 451
Princeton, NJ 08544

In astrophysics, plasmas occur under very extreme conditions. For example there are ultra strong magnetic fields in neutron stars, relativistic plasmas around black holes and in jets, extremely energetic particles such as cosmic rays in the interstellar medium, extremely dense plasmas in accretion disks, and extremely large magnetic Reynold's numbers in the interstellar medium. These extreme limits for astrophysical plasmas make plasma phenomena much simpler to analyze in astrophysics than in the laboratory. An understanding of such phenomena often results in an interesting way, by simply taking the extreme limiting case of a known plasma theory. I will describe one of the more exciting examples. I will attempt to convey the excitement I felt when I was first exposed to it. However, not all plasma astrophysical phenomena are so simple. There are certain important plasma phenomena in astrophysics, which have not been so easily resolved. In fact a resolution of them is blocking significant progress in astrophysical research. They have not yet yielded to attacks by theoretical astrophysicists nor to extensive numerical simulation. I will attempt to describe one of the more important of these plasma-astrophysical problems, and discuss why its resolution is so important to astrophysics. This significant example is fast, magnetic reconnection. Another significant example is the large-magnetic-Reynold's-number MHD dynamos.

MASTER

I. INTRODUCTION

Plasma physics plays an important role in a number of areas of astrophysics. I would like to discuss two such areas that I have been involved with for a long time. The first area is cosmic rays. In this field plasma physics has already made large contributions. There are many parallels with fusion physics.

The second area is magnetic reconnection. In this area the major understanding lies ahead, although considerable progress has been made.

I will survey these fields, describe some of my interactions with them and try to say what direction they are going.

II. COSMIC RAY RESEARCH

First Cosmic Ray research.

What are cosmic rays? Their early discovery was very exciting [1]. The first indication of their existence and the suspicion they were cosmic came from the simple electroscope. It was found that after the electroscope was charged and the gold leaves repelled each other, that the charge leaked off through finite conductivity of the air in the chamber. What causes this conductivity? Although it was Coulomb himself who first raised this question, it was not appreciated that ions were being produced in the air by energetic very penetrating particles. At the turn of this century radioactivity in the surroundings was suspected, but that this was not the complete story was shown by Hess. In a dangerous balloon flight to 16,000 feet he showed that the rate of ion production increased with height, proving that the cosmic rays came from above. As a result of his work and others, one can draw a plot of ion production versus depth below the top of the atmosphere. (See figure 1.)

Not much further progress occurred after this until the Geiger-Muller tube in 1929 made more precise measurements possible. During the thirties and forties the pattern of the interaction of the cosmic rays with air was unfolded as well as the existence of many new elementary particles. This was a truly heroic effort. The key results are shown in figure 2. This figure shows how the cosmic rays evolve through the atmosphere.

The nuclear particles (NA) are actually the cosmic particles raining on the earth at the rate of $1/\text{cm}^2 \text{ sec}$ summed over all directions. When they strike the atmosphere they penetrate only a short distance and then make a collision with the nucleus of a Nitrogen or Oxygen atom. In the resulting nuclear explosion pions are made. The neutral pion decays into photons initiating a cascade shower which at first grows and then decays. This is the E component. The charged pions decay into charged muons which penetrate all the way to the surface of the earth and are actually the main secondary particles at the ground. this is the M component. They produce the ions in found in the electroscope.

What are the primary particles? They consist mostly of protons, but there are heavier elements. They are all fully stripped. Figure 3 compares the abundance of the different species of cosmic rays which arrive at the earth with the accepted cosmic abundance of the elements. The relative abundance is very similar to the cosmic abundance with certain exceptions.

These exceptions are: an overabundance of the rare elements Li, Be, and B, and an underabundance of iron. These facts can be understood if it is postulated that the cosmic rays have been stored in the galaxy for millions of years. As they pass through the interstellar medium the more abundant elements are broken up by collisions with the interstellar medium atoms leading to the production of the light elements. How long are the cosmic rays stored in the galaxy? The assumption that the total number of Li, Be, and B nuclei in the cosmic rays are secondary and produced by spallation gives a lower limit to the age or storage time.

The existence of iron which can not survive too long a time against spallation gives an upper limit.

The best number for the age, in terms of gm/cm^2 , is $5 \text{ gm}/\text{cm}^2$ which converts to a time of 3 million years for a density of 1.0 interstellar medium atoms $/\text{cm}^3$ and 15 million years for $n = 0.1 \text{ atoms}/\text{cm}^3$. The presence of radioactive secondaries such as Be^{10} gives the absolute time as approximately 20 million years.

Thus, from the composition of cosmic rays one can state that they are stored in the galaxy for 20 million years, but because of the higher density in the galactic disk they spend only 3 million years in the disk and the remainder of their life in the lower density regions of the halo.

Other evidence indicates that cosmic rays fill the entire galaxy. This evidence is from the background synchrotron radio emission. Taking the fraction of one percent of the cosmic rays to be relativistic electrons and from the known value of the galactic field of 3×10^{-6} gauss, this map is consistent with cosmic rays filling the galaxy. From similar evidence other galaxies also are known to be filled with cosmic rays.

Other important properties of cosmic rays are their energy spectrum, and their anisotropy. The integral energy spectrum is remarkably close to a perfect power law up to about 10^{15} electron volts, as one can see by the plot in figure 4. They extend up to extremely high energies with of order joules per cosmic ray at the highest observed energies. Also, they are remarkably isotropic, better than one part in 10^{-3} over all directions.

Putting these facts together with what we know about the interstellar medium we arrive at the following picture. The galactic disk has a predominantly toroidal field of magnitude about 3 microgauss, whose topology is a little uncertain. (It is not known whether the field lines are entirely enclosed in the galactic disk are connected to the intergalactic medium.) In such a field the gyroradius of a cosmic rays is $10^{12} \epsilon_{\text{GeV}}$ centimeters, where ϵ_{GeV} is the energy

per nucleon of a cosmic rays in Gev. This may seem very large by terrestrial standards, but it is actually very tiny compared to the thickness of the galactic disk which is 10^{21} centimeters. The relativistic gyration period is $30\epsilon_{Gev}$ seconds.

Taking the scale of variation of the magnetic field to be 10^{21} centimeters, the thickness of the galactic disk, one finds that the perpendicular drift velocity of cosmic rays is 10 cm/s , an extremely small value. Thus, the cosmic rays are completely tied to the magnetic lines of force of the interstellar magnetic field (see figure 5). How they actually escape the galactic disk is a puzzle. From the very small anisotropy we know that their bulk motion parallel to the interstellar magnetic field is of order 200 km/sec. , so that in 3 million years they only go $600 \text{ parsecs} \approx 2000 \text{ lightyears}$. The reason for the small anisotropy is not Coulomb collisions with atoms in the interstellar medium, since the mean free path for Coulomb collisions at such high energies is 10^{12} lightyears. It is also not nuclear collisions. In fact, as will appear, it is by pitch angle scattering off of hydromagnetic waves that the isotropy is maintained.

If the interstellar magnetic field is open to the intergalactic medium, the small drift velocity confines the cosmic rays easily for 3 million years (figure 5a). If the interstellar magnetic field is closed in the galactic disk, the escape might be by the cosmic rays blowing bubbles in the interstellar magnetic field (figure 5b). (In fact, the cosmic ray pressure is comparable with the gas pressure and larger than the magnetic pressure.) Alternatively these bubbles may also be blown by a series of supernova blast waves. These bubbles must grow rapidly, otherwise the lifetime of the cosmic rays in the galactic disk is too long. In this picture the cosmic rays actually must spend five sixths of their lifetime in these bubbles continually returning to the disk from the bubble many times. From a simplified point of view the two pictures, showing open lines and closed lines, are quite similar.

Let me interrupt my discourse to tell a little personal history showing how I got involved with the problems of cosmic rays. Also, how I gradually came to an understanding of the

plasma instability which plays such a major role in the astrophysical behavior of cosmic rays.

I worked almost exclusively on controlled nuclear fusion problems up till 1966, when Ira Bernstein persuaded me to come to Yale University for a more academic career. When I got there I was somewhat at a loss as to what to work on, so I decided to see if there was any simple way to accelerate cosmic rays in the interstellar medium by using plasma physics. I embarked on this without doing much background reading, more for fun than for making a serious effort. This is often my *modus operandi*. I find out what need to know about a problem through conversations with my friends, rather than by reading. It leads to lots of mistakes, but I usually end up understanding the problem better.

Anyway, I started thinking about how to accelerate cosmic rays. For this electric fields are needed. There are no electric fields $\mathbf{E}$ parallel to the magnetic field $\mathbf{B}$, because of the high effective electrical conductivity of the interstellar medium. If $\mathbf{E}$ is perpendicular to $\mathbf{B}$, and constant, its accelerative power for a cosmic ray with velocity $\mathbf{v}$, $\mathbf{E} \cdot \mathbf{v}$ averages out. This is because the cosmic rays gyrate about the interstellar magnetic field once a minute. Thus, for effective acceleration $\mathbf{E}$ must change in such a way so that the cosmic rays see an electric field along $\mathbf{v}$ which does not average out. In other words, there must be a resonance with its cyclotron motion. If the electric field is produced by Alfvén waves this resonance condition is

$$\omega - \mathbf{k} \cdot \mathbf{v} = \pm \Omega \quad (1)$$

Now, $\omega \approx kV_A$ is small compared to $k v$ so the resonance condition reduces to

$$-k_{\parallel} v_{\parallel} \approx \pm \Omega \quad (2)$$

Thus, $k_{\parallel}$ must be of order $\Omega/v_{\parallel}$, approximately the reciprocal of the the cyclotron orbit of the cosmic ray. Here, $v_{\parallel}$ and $k_{\parallel}$ are the components of $\mathbf{v}$ and $\mathbf{k}$ parallel to $\mathbf{B}$. (A similar restriction holds for acceleration by electric fields carried by other waves.) Hence, I needed

to find an electric field that varies over the distance comparable to the gyroradius of the cosmic rays, $\approx 10^{12} \epsilon_{gev}$ cm.. Alfvén waves are the natural waves in the interstellar medium at this wave length.

Can they lead to acceleration? I argued as follows: If the Alfvén waves in the interstellar medium occur in wave packets that are close packed, then a given cosmic ray will pass $N = 10^{12}$ such wave packets. Some interactions of the cosmic ray with the wave are positive and some negative. An estimate of the net acceleration per wave packet is: its mean electric field E (without regard to sign) times the elementary charge times the length of the wave packet, reduced by the factor $\sqrt{N} \approx 10^6$. Even with this reduction, a total change in the energy of 1 Gev is produced during the lifetime of the cosmic rays of a million years, by an electric field of only 10^{-12} volt/cm. This means that Alfvén waves with very small relative amplitude $\delta B/B \approx 10^{-3}$ can produce enough acceleration to be interesting.

I mentioned this to a friend who pointed out an inconsistency in this estimate. If all the waves are going in the same direction, which has nothing to do with my estimate, then in their frame of motion $E = 0$, and there is no acceleration at all! Since $V_A \ll c$, one expects the total change in energy to be at most $(V_A/c)\epsilon$ no matter how many waves there are. In fact, there is only pitch angle scattering in the wave frame and the change that I estimated was indeed just due to pitch angle scattering in this frame. Since only a very small change in energy $\Delta\epsilon$ is associated with a finite change in the pitch angle, θ , pitch angle scattering must be much more efficient than energy scattering.

The error in my estimate of the change in energy arose because I kept only the random change in $\Delta\epsilon$ and forgot to include the systematic quadratic change in $\Delta\epsilon$. This second term mainly compensates for the random gain in energy. (The same mistake can, of course, be made in estimating the gain in electron energy from scattering off of heavy ions.) The secular quadratic term is positive if the particles meet the wave head on and negative if they

are going in the same direction as the wave and faster. Combining the terms representing the quadratic and the random effects one find that most of the acceleration disappears.

But this is logical. Pitch angle scattering is due to magnetic forces of magnitude $\mathbf{v} \times \mathbf{B}/c \approx B$ which are much larger than the electric forces E (by the factor c/V_A). Thus, although I had failed to find an acceleration mechanism for cosmic rays I had found a very strong scattering mechanism whose magnitude works out to be $\nu \approx \Omega(\delta B/B)^2$. Perhaps this was the solution of the trapping mechanism that could account for the isotropy of cosmic rays, and the finite but very long lifetime of the cosmic rays . [Because of my ignorance of the literature, I did not appreciate that this wave mechanism had been suggested 10 years earlier by Leverett Davis.]

Even though the wave amplitudes necessary to produce the required isotropy of the cosmic rays are small, there still has to be a source for the waves. What is it? What could produce Alfvén waves in this wave length range.

To solve this problem I again made a false start.

Many stars disturb the regions around themselves by stellar winds that extend over a distance comparable to the cosmic ray gyration radii. These stars must emit Alfvén waves by the Cerenkov mechanism. Could they emit enough waves for my purpose? I estimated the amount and found it was a little too small. When I presented these ideas in public Alar Toomre told me that I should talk to Don Wentzel. I did and he told me that Ian Lerche and he had done some calculations of the instability of Alfvén waves in the presence of cosmic ray anisotropy. This was a great idea! Their calculations were quite complicated because they were very general. I carried out similar calculations, but restricted myself to only parallel propagating waves. These calculations were much easier. The linearized result was very simple. If the cosmic rays tried to propagate with a mean bulk velocity $v_D > V_A$, the waves that cyclotron resonate with the cosmic rays become unstable. If the cosmic rays

were completely isotropic in a frame moving with the velocity v_D , the growth rate formula was even simpler. In fact, one can write the growth rate Γ of the Alfvén waves as

$$\Gamma = C \frac{\pi}{4} \Omega_0 \frac{v_D - V_A}{V_A} n_{cr}(\epsilon)/n \quad (3)$$

where Ω_0 is the nonrelativistic cyclotron frequency, $n_{cr}(\epsilon)$ is the number of cosmic rays with energy greater than ϵ , n is the density of the interstellar medium and C is a constant of order unity whose value depends on the energy spectrum of the cosmic rays [4].

Therefore, we may conclude that the origin of the waves is the cosmic rays themselves. If they try to escape from their source with bulk velocity greater than V_A , waves grow, and when they develop to a large enough amplitude they scatter the cosmic rays waves locking them in. Wentzel arrived at the same simplified results at the same time that I did and we published almost simultaneously. He beat me by a month [5,6].

III. THE COSMIC RAYS TRAP THEMSELVES//

Because of the simplicity of the result for the growth rate, it is easy to derive by a simple intuitive argument. The process is of course the analogy of Landau damping for plasma oscillation.

Consider a fixed small amplitude Alfvén wave packet of length comparable to the wave length λ of the waves in it (see figure 6). If a cosmic ray whose parallel gyroradius $r_{\parallel}$ is larger than λ , passes the wave, it hardly is disturbed by the wave (figure 6a). If $r_{\parallel}$ is much smaller than λ , then the cosmic ray follows the magnetic field adiabatically and again nothing happens (figure 6b). Consider a third cosmic ray for which $r_{\parallel}$ is comparable to λ (figure 6c). This cosmic ray tries to follow the magnetic field line to a certain extent. which is similar to the second cosmic ray, but it can't quite follow it, which is similar to the first cosmic ray. As a result, its pitch angle steps by about α , the angular deviation of the magnetic line of force from the direction of the mean field. Thus, the change in the

pitch angle is $\Delta\theta \approx \alpha \approx \delta B/B$. This change is random in sign. The cosmic ray encounters $v_{\parallel}/\lambda \approx k_{\parallel}v_{\parallel} \approx \Omega$ such wave packets per unit time. Adding the squares of $\Delta\theta$ we find that the cosmic ray pitch angle scattering formula is

$$\nu \equiv \frac{(\Delta\theta)^2}{t} \approx \Omega \left(\frac{\delta B}{B}\right)^2. \quad (4)$$

This is the same pitch angle formula that one arrives at by quasilinear theory.

What about the growth rate Γ ? Note the effect of the pitch angle scattering is that the cosmic rays are isotropized at the rate ν . Let the cosmic rays initially be isotropic in a frame moving with velocity $v_D > V_A$, (see figure 7). If we wait a time equal to the pitch angle scattering time $t \approx \nu^{-1} \approx B^2/\Omega(\delta B)^2$, the cosmic rays become isotropic in the wave frame with velocity V_A . They will have lost a total momentum equal to their (relativistic) mass times the difference in velocities $v_D - V_A$. This momentum goes into the wave momentum. Equating the rate of change of the wave momentum $\Gamma(\delta B)^2/4\pi V_A$, to the total change in the cosmic rays momentum divided by the scattering time t we get the formula of eq. (3) for the growth rate Γ .

This instability and the above formulas for the rate of pitch angle scattering and the wave growth rate are extremely important in cosmic rays research. They state that if the cosmic rays have a net streaming velocity in excess of the Alfvén speed, V_A , then Alfvén waves grow exponentially and pitch angle scatter the cosmic rays until their drift velocity is reduced to V_A . The typical growth time in the interstellar medium is relatively short, $\approx 10^4$ years, compared to their lifetime. The amplitude of the waves will be just such as to maintain a mean isotropy $\delta \approx V_A/c$ against whatever process is driving the cosmic rays towards anisotropy.

An example makes this clear. Let a source of cosmic rays be a distance L along a magnetic line from a sink. Then δ must be V_A/c . But $\delta \approx \lambda_{mfp}/L$ where λ_{mfp} is the mean free path for pitch angle scattering. From this we can estimate λ_{mfp} and thus $\nu \approx r_{\parallel}/\lambda_{mfp}$.

From this we can find $\delta B/B$ from eq.(4). Now the growth rate formula guaranties that this equilibrium situation will be reached if Γ is large enough that the waves can exponentiate during the time the source is on.

All these results refer to parallel propagating Alfvén waves. What about obliquely propagating magnetohydrodynamic waves? What about other waves? In fact, oblique magnetohydrodynamic waves have essentially the same effect, but reduced by Bessel function form factors. Also there is strong damping of these waves in high β plasmas. The same applies to other waves. The ion acoustic wave has a lower threshold for instability. if $c_s \ll v_A$, i.e. $v_D > c_s$. However, it is strongly damped by ion Landau damping and will not be excited. The plasma oscillations have such high phase velocities that there are only excited by Cerenkov resonance. They too are strongly damped by electron-ion collisions, even if their interaction with cosmic rays were unstable, which is generally not the case. The other modes such as whistler modes etc. have much higher phase velocities and thresholds. Only the parallel propagating Alfvén wave seems important, This is the one easiest to calculate.

What about linearized damping of parallel propagating Alfvén waves? Linearized damping in a fully ionized medium is extremely weak. However, the interstellar medium is only partially ionized over large regions. The situation is as follows:

Consider the motion of the plasma and the neutral medium separately. Only the ions experience magnetic forces. As they move they normally drag the neutrals along by collision. This is true if the wave frequencies are small compared to the collision frequency. But for the Alfvén waves of interest the Alfvén frequencies are much larger than the ion-neutral collision frequency. Thus, to first approximation the plasma moves independently of the neutrals. The wave velocity is really the Alfvén speed which includes only the plasma mass. This is considerably larger than the normal Alfvén speed which includes the total mass.

However, the ion-neutral collision rate ν_{in} can be larger than the growth rate Γ , so that

before the waves grow substantially a collision of an ion can occur with a neutral removing half the energy. This ambipolar effect, as it is called in astronomy, damps the wave at the rate $\nu_{in}/2$. Thus Γ must be larger than the damping rate, $\Gamma^* = \nu_{in}/2$, for instability. The stability condition thus allows faster drift velocities than the Alfvén speed. As a matter of fact, as one goes to high energy, the number of cosmic rays decreases and the self-trapping becomes weaker. The bulk drift velocity increases as $\epsilon^{3/2}$, and the lifetime decreases as $\epsilon^{-3/2}$. This was the situation in the late sixties. In the 70's it was discovered that large parts of the interstellar medium were hot and fully ionized. The linearized damping in these regions is negligible, and nonlinear damping becomes important instead. This leads to a lifetime proportional to $\epsilon^{-0.8}$, a result with an exponent about one half as large as that for partially ionized plasmas. The observed life time from the observation of the change in composition with energy is proportional to $\epsilon^{-0.6}$, so the agreement is reasonably good.

IV. THE ORIGIN OF COSMIC RAYS//

I should now like to mention the role that the cosmic ray instability has played in the theory of the origin of cosmic rays.

In the sixties and early seventies the favored origin for cosmic rays was supernovae. The large energy requirements for cosmic rays pointed to supernovae. However, if cosmic rays were created during the supernova explosion itself, they would become trapped by the instability as the supernova bubble grows, and they would lose a lot of energy adiabatically through the work they do in the expansion (see figure 8). Eventually, they escape along the magnetic field lines when the velocity of expansion drops to a value near the Alfvén velocity. The total linear expansion that occurs from the instant of the explosion itself until the release of the cosmic rays, is about a factor of five hundred. Adiabatic expansion reduces the energy of the cosmic rays by just this linear factor. Thus, adiabatic losses are extremely severe.

To produce the typical GeV cosmic rays observed at the earth, the energy that would have to be given to each of cosmic rays in supernovae must be five hundred GeV, an absurdly large result. The power requirements also become much larger; more than the total energy available in the supernova must go into cosmic rays. The question of whether the cosmic rays are actually trapped along the field lines extending out of the supernova bubble, was hotly debated during the seventies. I remember one or two exciting public arguments with Sterling Colgate on this matter. He claimed that in nuclear fusion devices it is impossible to confine energy, and the same must be true in supernova remnants. Anyway, after all the results were in, self trapping won. The growth rates were fast enough to overcome even nonlinear damping the cosmic rays must indeed be trapped.

By the end of the seventies, new ideas concerning the origin of cosmic rays emerged. Blandford and Ostriker realized that after even after the supernova bubble had become large, the supernova energy was still present in the form of kinetic energy of the expanding bubble. This energy could be tapped for cosmic rays acceleration, but at a time such that the bubble was so large that there need not be serious adiabatic loss. They invented a very clever method of acceleration which is natural, unavoidable, and produces the correct energy spectrum [7].

In this method the plasma instability again enters. It proceeds as follows: (see figure 9).

Imagine that there are cosmic rays present in the region upstream from the shock. Let $v_s \gg v_A$. Upstream we start off with the bulk velocity of the cosmic rays $v_D \approx v_s$ in the shock frame. As the fluid crosses the shock the medium is slowed down to a speed $v_s/4$. The cosmic rays are unaffected by the shock itself, and therefore have $v_D \approx v_s$ downstream. They thus have a bulk velocity relative to the plasma of $3v_s/4$ which is large enough to drive the plasma instability, the waves grow to significant amplitude, and self trapping sets

in. The cosmic rays are scattered back across the shock into the upstream medium where the bulk velocity v_D becomes reduced below v_s , causing wave amplification upstream. The situation then settles down to a steady state and the cosmic rays are decelerated upstream by cosmic waves.

Now, because the cosmic rays have actual velocities of c , but the bulk velocity of $\approx v_s$, there is of order c/v_s pitch angle scatterings in the neighborhood of the shock. This means on the average every cosmic rays crosses the shock c/v_s times.

But every time the cosmic ray crosses the shock it changes its energy by scattering off of the moving medium. Upstream, in the frame of the shock, it gains energy $(v_s/c)\epsilon$, and downstream it loses energy $1/4(v_s/c)\epsilon$, so that the net change per double crossing is $3/4(v_s/c)\epsilon$, which is always positive. (This is first order Fermi acceleration.) Now, in order that the bulk velocity of the cosmic rays be reduced from upstream to downstream, enough scattering must occur so that a typical cosmic ray crosses the shock of order c/v_s times, thus leading to a gain of energy ϵ on the average. But some cosmic rays gain more energy, and some less energy, by the random statistical nature of the process. The final answer is that the cosmic rays come out of the shock downstream with a power law distribution in energies proportional to ϵ^{-2} . The spectrum is a little less steep if the shock wave is not infinitely strong, that is if v_s is not very large compared with v_A . The observed spectra distribution is roughly proportional to $\epsilon^{-2.5}$. This very natural process is now the leading candidate for the origin of cosmic rays.

I hope that I have convinced you that plasma physics employing standard techniques has already made an important contribution to the understanding of cosmic rays in their astrophysical context.

Now, I turn to the less well understood subject of magnetic reconnection, for which the understanding of the basic plasma processes still lies in the future.

V.MAGNETIC RECONNECTION//

A second major area in which advanced plasma ideas impinge on astrophysics is magnetic reconnection. This is a subject I have also played with off and on throughout my career. Let me summarize the situation.

First, what is magnetic reconnection? For an infinitely conducting plasma it is not supposed to happen. But, a plasma is really not infinitely conducting. Astrophysical plasmas are no more highly conducting than laboratory plasmas. Astrophysical plasmas have an effectively high conductivity because of their large size. The important quantification of the plasma conductivity is the skin time which is

$$t_{skin} = \frac{4\pi L^2}{\eta} \approx 10^{-7} L^2 T_{ev}^{3/2} \quad (5)$$

in cgs units. T_e is the electron temperature in electron volts. As an example consider the magnetic field in the galaxy. If we take L the thickness of the galactic disk $\approx 10^{21}$ cm, and $T_{ev} \approx 1$ then $t_{skin} \approx 10^{35}$ sec = $10^{27.5}$ years. On the other hand, experience has taught us that in three dimensional equilibrium one often has surface currents, thin regions which carry a finite amount of current, and over which the field changes abruptly.

Generally, these regions are not infinitely thin, but their thickness gets smaller and smaller as time progresses. Eventually the skin time, being proportional to the thickness squared, becomes sufficiently short that resistance becomes important. One loses the infinite conductivity assumption, and thus also flux freezing over these regions. Of course, the infinite conductivity assumption remains valid outside of these tiny regions. Why is the breakdown of flux conservation over tiny regions of importance? It is because many lines are crushed together in these regions, and lines passing through them extend to different areas far apart in space.

Thus, two lines of force whose ends are far apart become cut from themselves and

attached to each other. (See figure 10) This changes the topology and the dynamics of the system in a manner totally unexpected in infinite conductivity systems.

It is known that magnetohydrodynamic equilibria are really minimum energy states of the system, subject to boundary conditions and to topology of the magnetic field lines. A small change in topology can lead to a state which may have a much lower energy. Perhaps, this can be made clear by a hypothetical example from the solar corona due to Sweet [8].

Consider two pairs of sunspots, and let the pairs approach each other. Figure 11a and 11b show what happens if the corona were a vacuum. (Lines are labeled by their foot points.) But if one considers the electric field at N, one sees that the changing flux between N and the solar surface induces an electric field at N. If sufficient plasma were present, then large currents flow near N and one gets the situation of figure 11c which corresponds to flux freezing. Now, a narrow layer appears at N. If the two lines labeled 2 are reconnected, then the two lines labeled 3 approach N and are also reconnected. In this way we can get back to figure 11b. But notice that the two parts of line 3 in figure 11a are connected to the same pair of sun spots at their ends. After reconnection the two parts of 3 are effectively cut and resewn so that we get two new lines each passing from one pair of spots to the other. One can see almost intuitively that configuration of figure 11b has a lower energy than configuration of figure 11c. Thus, assuming large conductivity a great deal of work is done on the field by moving the sunspots, thus greatly increasing the energy. Then, at some point reconnection sets in taking the field to the lower energy state. The difference in the stored energy appears as heat, mechanical motion, radiation and acceleration. It was supposed that some such sequence of events underlies the solar flare phenomenon.

This picture was put forward by Sweet in the late 50's [8]. What was missing (and perhaps still) is a realistic estimate of the time scales.

Because of the complexity of the event, most calculations have been two dimensional.

The basic physics of reconnection involves two ideas (figure 12). First, the thinner the layer, the faster the oppositely directed fields will merge. Second, to counter this one also notes that the thinner the layer, the harder it is to remove the plasma on the incoming reconnecting lines from the reconnecting region. Thus, a compromise is reached. The thickness of the layer, δ cannot be too thin, else one can't get the plasma out of the way fast enough. Since the fastest velocity at which plasma can flow is the Alfvén speed, $\rho\delta v_A$ is the rate of removal of plasma. δ cannot be too thick, otherwise the reconnection velocity $v_{rec} = \eta/\eta$ brings the plasma together too slowly. Thus a compromise must be reached at $\delta^2 = L\eta/v_A$.

If we define t_R as t_{skin} for lengths comparable to the outside length, L supposedly the length of the layer, then the reconnection velocity is of order

$$v_{rec} \approx \sqrt{\frac{v_A L}{t_R}} \quad (6)$$

and the time to reconnect a finite amount of flux is the geometric mean of the dynamical time $t_D = L/v_A$ and t_R [8,9]. Thus,

$$t_{SP} = L/v_{rec} \sqrt{t_R t_D} \quad (7)$$

(SP stands for Sweet-Parker whose papers first laid this picture out clearly).

This time is easily estimated, and is found to be much longer ($\sim$ a day) than the energy release time of a solar flare. Thus, although up to this point in our discussion the sequence of events might be said to compellingly lead to an explosion, it turns out that the rates are far too slow.

From this point on major attention was concentrated on the physics in the current layer. How could it be made to go faster? In 1963 Petschek introduced a new idea, which captured everybody's imagination [10]. Instead of a plane slab for the current layer (figure 12), for which the removal of the plasma was such a difficulty, maybe a funnel geometry is appropriate (figure 13).

Only a part of the line need be merged. Once a part of the line is merged, the tension present in the field can lift that part of the plasma away, which is not actually passing through the Sweet-Parker reconnection region. Thus, the problem of the removal of the plasma was greatly relieved. The Sweet-Parker layer could be much shorter than L , the macroscopic dimension. Petschek got an effectively complete although qualitative dynamical solution of the problem. He showed that almost any length of the Sweet-Parker layer was possible; not just the Sweet-Parker reconnection rate for a layer of length L . To make his theory complete Petschek then showed that there was an absolute upper limit of order $v_A/\ln(t_R/t_A)$ for the reconnection velocity. Now this absolute limit was seized upon by almost the entire astrophysical community and given the name "Petschek reconnection rate". There was an intense debate over whether reconnection proceeded at the Sweet-Parker rate or this Petschek rate. Somehow, most of Petschek's arguments were forgotten.

Now, one can make a qualitative estimate of the lengths in the Petschek picture, that would occur in the solar corona. For this rapid reconnection, the length of the Sweet-Parker layer must be less than 1 cm and the thickness about a tenth of this!

Let us go back to our original picture of the two pairs of sunspots coming together (figure 14), and try to follow what might happen in time. As the sunspot pairs move the equilibrium situation is as in figure 14a. A current layer forms near N . The layer should get thinner and thinner at some rate, which is somehow related to the rate at which the sunspot pairs approach each other. As the layer gets thinner reconnection starts at first very slowly. But there comes a time when the reconnection rate is as fast as the rate of thinning of the layer. This actually comes long before the time which the layer becomes thin enough to reach the so-called Petschek rate. I call it so-called because remember Petschek actually said that reconnection should be able to go at any rate. It seems to proceed as fast as it needs to, to match the changing boundary conditions. But the reconnection itself prevents the layer from becoming short enough for really rapid reconnection.

Syrovatskii approached the problem somewhat differently, but in such a manner that the outside boundary is properly included [11]. Let us seek a magnetostatic solution of the field in figure 14a consistent with the topology of the lines. As we have learned from experience, such an equilibrium demands a surface current. Treat this surface current as infinitely thin. One gets the field pattern of figure 14a where everything is totally determined by the position of the foot points at the ends of each line. Now suppose we allow reconnection. The Petschek funnel is now outside of the current layer. The dynamics in the layer are those of Sweet-Parker. The Sweet-Parker reconnection rate changes the topology. Lines 2 reconnect and get out of the way at the Sweet-Parker rate, and more lines move in. If one stops moving the sunspot pairs, one freezes the position of the foot points and fixes the boundary conditions. Then reconnection proceeds, changing the topology at the Sweet-Parker rate. This rate is determined by two parameters: the length of the layer, and the magnetic field just outside of N. after a short time the line configuration is as in figure 14b. As reconnection proceeds these lengths and field strengths change in a manner determined by our mathematical procedure of solving for equilibria with given topology and boundary conditions. This procedure seems quite straightforward, and leaves no place for rapid reconnection on the Petschek time scale which requires a much shorter reconnection layer.

Now, in the mid 1980's, Biskamp did a number of remarkable numerical simulations of two dimensional reconnection, and found exactly these results [12]. The only difference from the above model is that he did not evolve through a system of static magnetohydrodynamic equilibria, but he included the full magnetohydrodynamics.

He simulated seven cases for different values of the plasma conductivity and different boundary conditions leading to different reconnection rates. A sketch of the field lines in two of the cases is shown in figure 15. In each case Biskamp varied the pressure on the boundary both of the incoming and outgoing flow to reach a steady state with the desired reconnection velocity. He found that in every case a current layer formed whose length and

current strength (equal to the piled up field outside the layer) depended on specific conditions and varied considerably with each of the seven cases. Again the Petschek funnel is evident, but it is not in the boundary layer. Its main effect is to shorten the current layer.

The reconnection rate is closely given by the Sweet-Parker rate with these local parameters in all seven cases. Biskamp determined a steady state solution, so the reconnection rate turned out to be just fast enough to accommodate the rate at which he introduced flux. I think that Biskamp's calculations lay to rest the controversy over whether reconnection is Sweet-Parker or Petschek. He finds that the reconnection rates are certainly faster than estimated by Sweet-Parker because the current layer is shorter than L , and the magnetic field is stronger than the field upstream. Magnetic flux has "piled up" leading to this stronger field. The reconnection rates are much slower than the maximum rate generally quoted as the Petschek rate, but they are quite consistent with Petschek's statement that reconnection can go at nearly any rate.

Reconnection seems to always take place just as fast as it has to, to be consistent with changing outside conditions. But this rate is probably only faster by a finite factor than the Sweet-Parker original estimate. In the configuration of figure 14, the original estimate of the time to reconnect a finite amount of flux was one day. By shortening the current layer it can be reduced to perhaps hours.

What happened to the solar flare? How can it release energy so much faster than plausible estimates lead to? The problem, I think, lies in the two-dimensional assumption. It is almost impossible to be certain about anything in fully three dimensional magnetohydrodynamics without difficult numerical computations. But nature is three dimensional. The sunspots shown in figure 14 are round, and extend no further out of the plane figure than they do in the plane of the figure. I would suppose that if one attempted to find a three dimensional equilibrium with given topology, that there would be many surface current lay-

ers. Further, the foot point motions in the two dimensional model have been hypothesized to be quite regular. Motion of foot points in the photosphere are observed to be much more complex. Suppose there were N surfaces each of length D out of the paper. Then let Sweet-Parker reconnection go on in each of them. Then the effective velocity of reconnection of magnetic field lines is of order $NDv_A\sqrt{t_A/t_R}$. In other words, let the reconnection be multiprocessed. Then it is conceivable that the required amount of flux for a flare can be reconnected in a much shorter time than the Sweet-Parker estimate.

I have no idea if this picture is correct. However, there is an observational result which seems to bear it out. Flares occur with all sorts of energies. However, the "big" ones are rarer than the small ones. In figure 16, the log of the number of flares with peak energy release greater than ϵ , $N(\epsilon)$, is plotted versus $\log(\epsilon)$ one gets a straight line of slope -1.8 with fairly high accuracy [13].

This is a strong indication that flares are statistical. To me this says that the flare is the sum of a lot of small reconnections each going off nearly simultaneously at the Sweet-Parker rate. Various attempts to model the flare as a cascade of reconnections do actually lead to such a power law. For example, this is done in the "avalanche" theory of flares of Lu and Hamilton [13]. To my taste this is a very satisfactory situation if true. It involves plasma physics that we are fairly certain of, and it avoids the nasty extrapolation to fast reconnection, which relies on such absurdly small lengths. However, it lifts the problem to a different and higher degree of difficulty, namely the study of complex three dimensional equilibria with complex topologies. Such a picture has been advocated by Parker for a number of years to explain coronal heating. It seems likely that it also underlies the solar flare phenomenon.

There are a number of much more extreme astrophysical cases in which a basic understanding of reconnection is important. It seems to me that when one goes to galactic

length scales, such a story is essential if substantial reconnection occurs. The extremely long Sweet-Parker time for a single reconnection is of order of 10^{17} years. I believe that it could be greatly reduced by having a large number of reconnecting regions acting simultaneously. Petschek's fast reconnection at the Alfvén speed is a chimera that has misled astrophysicists for years. It seems more reasonable to accept the Sweet-Parker reconnection rate that has been verified a number of times by numerical calculations and laboratory experiments.

The two main points are these: Petschek has shown that Sweet-Parker reconnection can go faster, as has been confirmed by Biskamp. The three dimensional nature of reconnection should lead to a multitude of reconnecting regions all acting at the slower rate but simultaneously.

Finally, following the ideas of Syrovatskii, it may be appropriate to treat the problem as an evolution through equilibrium states each depending on the boundary conditions and the topology of the magnetic lines and each containing many surface currents. This ignores the dynamics outside of the the layer. This is roughly justified if the velocities are subalfvenic outside of the layers. Each surface changes topology by reconnection through itself at the local Sweet-Parker rate.

VI. CONCLUSION//

I have tried to illustrate by two examples how plasma physics impinges on astrophysics in an important way. I chose these examples from my experience, because I was most familiar with them. Cosmic ray research has benefited greatly by the discovery of a normal plasma instability. Normal plasma techniques have unfolded most of its manifestations. The second example, magnetic reconnection, is much harder and comparably little effort has been devoted to understanding it. Success has been poor because very naturally two dimensional techniques have been applied to an essentially three dimensional problem. The Petschek upper limit which has inappropriately been denoted as the Petschek reconnection rate seems

of little relevance. The reconnection rate is more probably determined by the length of the layer and the pile up of the magnetic flux outside. These are determined by solving the problem outside of the current layer so that there is no single reconnection rate.

I anticipate that fresh work on this second problem by plasma physicists should have as big an impact on astrophysics, just as the work on cosmic rays research has done.

ACKNOWLEDGMENT

It is a pleasure to thank the referee for his careful reading of the manuscript and the many helpful changes which he suggested. This paper was prepared under the auspices of the department of energy contract No. DE-A-76-CHO-3073, the NSF Grant No. AST-9121847 and NASA Grant NAGW-2419.

REFERENCES

1. B. Rossi, *Cosmic Rays* (McGraw Hill, New York, 1964).
2. N. Lund *Cosmic Radiation in Contemporary Astrophysics* (edited by M. M. Shapiro, D Reidel Publishing Co., Dordrecht) p.1
3. Y. Fujimoto, H. Hasegawa, and M. Taketani, Supplement of the Progress of Theoretical Physics 30,32 , (1964)
4. R. M. Kulsrud and C. J. Cesarsky, Astrophysical Letters 8, 189 (1971).

FIGURE CAPTIONS

1. Intensity of cosmic rays as a function of atmospheric depth. The atmospheric depth plotted on the horizontal axis is the mass per unit area below the top of the atmosphere. The vertical scale gives the number of ion pairs per second per cm^2 in the detector. The figure is taken from reference 1.

2. Vertical intensity of three components of local cosmic rays as a function of atmospheric depth in grams per square centimeter. Curve NA represents nuclear-active particles with energies greater than about 1 Gev. Curve E represents electrons with energies greater than about 100 Mev. Curve M represents muon with energies greater than about 200 Mev. The figure is taken from reference 1.

3. Elemental abundances of arriving cosmic rays (line) compared with local galactic abundances (bars). The normalization of the vertical axis is to hydrogen. The figure is taken from reference 2.

4. Integral energy spectrum of primary cosmic rays. The figure is taken from reference 3.

5a. Magnetic field lines in the galactic disk-open model. The thickness of the disk is about 10^{21} centimeters. A cosmic ray is shown spiraling on one of the lines. the part of the line outside of the actual disk is dotted.

5b. Magnetic field lines in the galactic disk-closed model. the loop of the line on the left is a part blown out of the disc by supernovae explosions. The detached circle on the left represents this loop if the line can be detached by magnetic reconnection.

6. Interaction of a cosmic ray with an Alfvén wave in three cases.

6a. The pitch of the cosmic ray helix is long compared to the wave length and $\delta\theta = 0$.

6b. The pitch of the cosmic ray helix is short compared to the wave length and $\delta\theta = 0$.

6c. The pitch of the cosmic ray helix is comparable to the wave length and $\delta\theta \approx \alpha$, where α is the angle the line makes with the direction of the

undisturbed magnetic field.

7. Contours of the cosmic rays ray distribution function in momentum before and after pitch angle scattering. The first contour is a circle centered about the point at $p_{\parallel} = mv_D$ and the second is a circle centered about the point $p_{\parallel} = mv_A$.

8. A expanding supernova bubble. The cosmic rays are trapped in the expanding gas and lose energy adiabatically. The plasma instability keeps the cosmic rays from flowing freely along the line of force out of the bubble, since the bubble expands much faster than the bulk flow velocity along the field which is of order the Alfvén velocity.

9. Cosmic ray acceleration by a shock wave. In the frame of the shock the fluid velocity changes abruptly from v_s to $v_s/4$, so in the absence of the Alfvén wave instability, the bulk velocity of the cosmic rays relative to the fluid abruptly increases. The cosmic rays must cross the shock c/v_s times to reduce the relative bulk velocity from its value upstream to a value downstream consistent with the relative velocity being less than v_A .

10. Reconnection of lines of force in just a small region can change the topology of the field so that points far away that are not connected by field lines before reconnection become connected by field lines.

11. The lines of force from two pairs of sunspot on the solar surface labeled by the amount of flux between them and the solar surface. The pairs of sunspots are imagined to move towards each other.

11a. The field configuration before the sunspots move. The field is initially a vacuum field.

11b. The field configuration after the sunspots have moved in the case where the solar region is a vacuum field. Lines with the same flux underneath in the two diagrams have the same numbers.

11c. The field configuration after the sunspots have moved in the case where the region is a conducting plasma. These figures are taken from reference 8.

12. An enlargement of the region about the neutral point in figure 11c,

where reconnection of the lines is occurring. The plasma flows vertically out of this region at the alfvén speed. This limits the thickness of the region as well as the rate at which the lines are brought in to reconnect. This figure is taken from reference 8.

13. If tension forces are taken into account as suggested by Petschek the Sweet-Parker reconnection layer can be made shorter and reconnection can proceed faster. The line configuration for this to happen is shown in this figure. This figure is taken from reference 10.

14. The field configurations at two stages during reconnection. The sunspots pairs are fixed. The field pattern does not change substantially but the line labels do. The manner of connection of the lines to the sunspot pairs defines the topology. It is proposed that the equilibrium of the lines, the length of the Sweet-Parker layer and the field strength just outside of the layer are determined in terms of this topology. As the topology changes due to reconnection these parameters change. Conversely, the rate of reconnection is given in terms of the length of the Sweet-Parker layer and the strength of the field just outside of it. This figure is taken from reference 8.

15. A sketch of the field line pattern in two of Biskamp's numerical simulations of magnetic reconnection [12]. Figure A is for low conductivity and figure B is for high conductivity. The rate of reconnection is the same in the two cases, but the length and strength of the current layer are different to compensate for the different conductivities.

16. Distribution of flare occurrence rate versus peak count rate for flares. The distribution has a logarithmic slope of -1.8 . This figure is taken from reference 13.

PPPL3 #94X0280

Fig. 1

PPPL#94X0281

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11

PPPL#94X0291

Fig. 12

PPPL#94X0292

Fig. 13

Fig. 14

Fig. 15

PPPL#94X0294

Fig. 16

EXTERNAL DISTRIBUTION IN ADDITION TO UC-420

Dr. F. Paoloni, Univ. of Wollongong, AUSTRALIA	Dr. H. Yamato, Toshiba Res & Devel Center, JAPAN
Prof. R.C. Cross, Univ. of Sydney, AUSTRALIA	Prof. I. Kawakami, Hiroshima Univ., JAPAN
Plasma Research Lab., Australian Nat. Univ., AUSTRALIA	Prof. K. Nishikawa, Hiroshima Univ., JAPAN
Prof. I.R. Jones, Flinders Univ, AUSTRALIA	Librarian, Naka Fusion Research Establishment, JAERI, JAPAN
Prof. F. Cap, Inst. for Theoretical Physics, AUSTRIA	Director, Japan Atomic Energy Research Inst., JAPAN
Prof. M. Heindler, Institut für Theoretische Physik, AUSTRIA	Prof. S. Itoh, Kyushu Univ., JAPAN
Prof. M. Goossens, Astronomisch Instituut, BELGIUM	Research Info. Ctr., National Instit. for Fusion Science, JAPAN
Ecole Royale Militaire, Lab. de Phy. Plasmas, BELGIUM	Prof. S. Tanaka, Kyoto Univ., JAPAN
Commission-European, DG. XII-Fusion Prog., BELGIUM	Library, Kyoto Univ., JAPAN
Prof. R. Bouciqué, Rijksuniversiteit Gent, BELGIUM	Prof. N. Inoue, Univ. of Tokyo, JAPAN
Dr. P.H. Sakanaka, Instituto Fisica, BRAZIL	Secretary, Plasma Section, Electrotechnical Lab., JAPAN
Prof. Dr. I.C. Nascimento, Instituto Fisica, Sao Paulo, BRAZIL	Dr. O. Mitarai, Kumamoto Inst. of Technology, JAPAN
Instituto Nacional De Pesquisas Espaciais-INPE, BRAZIL	Dr. G.S. Lee, Korea Basic Sci. Ctr., KOREA
Documents Office, Atomic Energy of Canada Ltd., CANADA	J. Hyeon-Sook, Korea Atomic Energy Research Inst., KOREA
Ms. M. Morin, CCFM/Tokamak de Varennes, CANADA	D.I. Choi, The Korea Adv. Inst. of Sci. & Tech., KOREA
Dr. M.P. Bachynski, MPB Technologies, Inc., CANADA	Prof. B.S. Liley, Univ. of Waikato, NEW ZEALAND
Dr. H.M. Skarsgard, Univ. of Saskatchewan, CANADA	Inst of Physics, Chinese Acad Sci PEOPLE'S REP. OF CHINA
Prof. J. Teichmann, Univ. of Montreal, CANADA	Library, Inst. of Plasma Physics, PEOPLE'S REP. OF CHINA
Prof. S.R. Sreenivasan, Univ. of Calgary, CANADA	Tsinghua Univ. Library, PEOPLE'S REPUBLIC OF CHINA
Prof. T.W. Johnston, INRS-Energie, CANADA	Z. Li, S.W. Inst Physics, PEOPLE'S REPUBLIC OF CHINA
Dr. R. Bolton, Centre canadien de fusion magnétique, CANADA	Prof. J.A.C. Cabral, Instituto Superior Tecnico, PORTUGAL
Dr. C.R. James,, Univ. of Alberta, CANADA	Prof. M.A. Hellberg, Univ. of Natal, S. AFRICA
Dr. P. Lukác, Komenského Univerzita, CZECHO-SLOVAKIA	Prof. D.E. Kim, Pohang Inst. of Sci. & Tech., SO. KOREA
The Librarian, Culham Laboratory, ENGLAND	Prof. C.I.E.M.A.T, Fusion Division Library, SPAIN
Library, R61, Rutherford Appleton Laboratory, ENGLAND	Dr. L. Stenflo, Univ. of UMEA, SWEDEN
Mrs. S.A. Hutchinson, JET Library, ENGLAND	Library, Royal Inst. of Technology, SWEDEN
Dr. S.C. Sharma, Univ. of South Pacific, FIJI ISLANDS	Prof. H. Wilhelmson, Chalmers Univ. of Tech., SWEDEN
P. Mähönen, Univ. of Helsinki, FINLAND	Centre Phys. Des Plasmas, Ecole Polytech, SWITZERLAND
Prof. M.N. Bussac, Ecole Polytechnique,, FRANCE	Bibliotheek, Inst. Voor Plasma-Fysica, THE NETHERLANDS
C. Mouttet, Lab. de Physique des Milieux Ionisés, FRANCE	Asst. Prof. Dr. S. Cakir, Middle East Tech. Univ., TURKEY
J. Radet, CEN/CADARACHE - Bat 506, FRANCE	Dr. V.A. Glukhikh, Sci. Res. Inst. Electrophys. Apparatus, USSR
Prof. E. Economou, Univ. of Crete, GREECE	Dr. D.D. Ryutov, Siberian Branch of Academy of Sci., USSR
Ms. C. Rinni, Univ. of Ioannina, GREECE	Dr. G.A. Eliseev, I.V. Kurchatov Inst., USSR
Preprint Library, Hungarian Academy of Sci., HUNGARY	Librarian, The Ukr.SSR Academy of Sciences, USSR
Dr. B. DasGupta, Saha Inst. of Nuclear Physics, INDIA	Dr. L.M. Kovrizhnykh, Inst. of General Physics, USSR
Dr. P. Kaw, Inst. for Plasma Research, INDIA	Kernforschungsanlage GmbH, Zentralbibliothek, W. GERMANY
Dr. P. Rosenau, Israel Inst. of Technology, ISRAEL	Bibliothek, Inst. Für Plasmaforschung, W. GERMANY
Librarian, International Center for Theo Physics, ITALY	Prof. K. Schindler, Ruhr-Universität Bochum, W. GERMANY
Miss C. De Palo, Associazione EURATOM-ENEA , ITALY	Dr. F. Wagner, (ASDEX), Max-Planck-Institut, W. GERMANY
Dr. G. Grosso, Istituto di Fisica del Plasma, ITALY	Librarian, Max-Planck-Institut, W. GERMANY
Prof. G. Rostangni, Istituto Gas Ionizzati Del Cnr, ITALY	