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PROBLEMS RELATED TO MACRO-SCOPIC ELECTRIC FIELDS IN THE MAGNETOSPHERE

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PROBLEMS RELATED TO MACROSCOPIC ELECTRIC FIELDS IN THE MAGNETOSPHERE

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1 Introduction

In the early stages of magnetospheric research, the electric field was generally regarded as rather unimportant. This was perhaps natural at a time when low energy particles could not be measured and when, for lack of observational data, it was generally believed that the whole magnetosphere was pervaded by a "cool background plasma" with all the desirable properties of an ideal magnetohydrodynamic medium. Furthermore, the task of measuring the electric field presented comparatively great technical difficulties, which for some parts of the magnetosphere are not entirely overcome yet.

Although electric field measurements have now been conducted in the ionosphere for about a decade, important parts of the magnetosphere are still entirely virgin ground for direct electric field measurements.

Thus the electric field is still a kind of last frontier in magnetospheric exploration. Very recently important discoveries were made by the first direct electric field measurements at high altitude in auroral and polar regions.

It has also become apparent that the theoretical problems related to magnetospheric electric fields are no less important and demanding. The image of magnetospheric plasma as a simple magnetohydrodynamic medium that can be described by highly idealized models has been shattered. In its place emerges an image of great complexity. Rather than being a virtually resistanceless conductor of electric current along magnetic field lines, the real magnetospheric plasma appears capable of supporting strong magnetic-field-aligned of supporting strong magnetic-field-aligned electric fields, and corresponding voltage drops of many kilovolts. Mechanisms that may be responsible for this have been identified but are not readily amenable to satisfactory theoretical analysis. Direct observational data on electric fields in the outer magnetosphere are greatly needed.

2 Sources of magnetospheric electric fields

The large-scale electric field in the magnetosphere originates from internal as well as external sources. The internal sources are of basically two kinds: the "unipolar dynamo" constituted by the rotation of the ionosphere as a whole, and dynamo effects due to ionospheric winds. (In addition, the ionosphere can play an important role in redistributing electric fields of external origin.) The primary external source is of course the solar wind dynamo.

2.1 Internal sources

Of the internal sources, only the unipolar dynamo will be discussed in the present context.

In a nonrotating frame of reference, the Earth's ionosphere presents an electromotive force of nearly 100 kV to the external plasma. Most of the rotational e.m.f. is applied to the collision-dominated plasma inside the plasmapause, where the magnetic field lines are equipotentials, or nearly so. In this region the applied e.m.f. causes the plasma to corotate but has no particularly interesting consequences.

Outside the plasmapause the plasma is largely collisionless, and its response to electric fields is much less simple. This is in particular true for the plasma in magnetic flux tubes emanating from the auroral oval and the polar caps. It is worth noting that the part of the unipolar dynamo e.m.f. applied over this region is as large as about 10 kV.

The possibility that the terrestrial unipolar dynamo might be of interest in magnetospheric physics was recognized long ago by McIlwain (1969). Recently it has been invoked more specifically, e.g. in explaining the correlation between polar geomagnetic variations and the polarity of the interplanetary magnetic field (Volland, 1975a).

If the electrical coupling were sufficiently good, the high-latitude unipolar-dynamo field would make itself felt deep in the magnetosphere tail, although modified due to the inertia of the plasma throughout the high-latitude flux tubes. On the other hand, the electric field imposed from below may be more or less disconnected from the plasma in the more distant parts of the magnetic flux tubes. This requires substantial magnetic-field-aligned electric fields to exist in the high-latitude topside plasma (cf §3.2). It is conceivable that a plasma-dynamical mismatch necessitating decoupling is one of the driving forces for establishing parallel electric fields.

Together with the e.m.f. imposed from the solar wind the unipolar dynamo creates an electric field distribution which is asymmetric in a way that depends on the direction of the solar wind magnetic (and hence electric) field. This asymmetric field distribution was first observed by Heppner (1972). The actual field is of course not given by a simple superposition of the external and internal electromotion forces but is distributed by the complex response of the plasma, which depends not only on its "conductivity" but also on the dynamics of its motion (cf. §3).

The solar-wind dynamo

The energy flux carried by the solar wind over an area such as the cross-section of the magnetosphere would be amply sufficient to continuously support strong geomagnetic activity. Yet, it is only occasionally that any appreciable fraction of this power is admitted into the magnetosphere. The problems of how the solar wind interacts with the magnetosphere are not discussed in the present paper. Here we may just notice that the solar wind can be looked upon as an MHD dynamo whose e.m.f. has a magnitude of the order of 10^5 V, but applied at various angles relative to the magnetosphere. In terms of energy extraction from this dynamo, it is interesting to note the observational result that the magnetosphere acts like linear resistor in series with a rectifier that has a dawn-to-dusk conduction direction. This kind of behaviour manifests itself in a very clear way in the recent results of Burton et al. (1975). Plotted as a function of the east-to-west interplanetary electric field, E_y, the rate of energy injection into the magnetosphere is essentially zero for negative E_y and is proportional to E_y for $E_y>0$.

3 Electric fields and plasma convection

In studies of magnetospheric plasma convection, the concept of frozen-in magnetic field lines has been a favourite tool. This choice is understandable because of the elegance and convenience of the concept. As long as the magnetosphere was believed to be filled with a cool plasma having essentially infinite conductivity along magnetic field lines it was also very natural to use the frozen-in concept. Now the situation is different. In large regions of the magnetosphere the plasma has characteristic particle energies of the order of kiloelectronvolts, and electrical properties that are far from simple. In fact, there need not even exist a local relation between electric field and current, which is required for a relation such as the generalized Ohm's law to hold. There has, however, been widespread reluctance to abandon the frozen-field line concept, which has caused Alfvén, who introduced the concept of frozen field lines, to warn strongly against its use in low-density space plasma (Alfvén, 1968, 1976).

3.1 Energy-dependent differential convection

Even disregarding the possible field line unfreezing by magnetic-field-aligned electric fields, we may notice that already the finite particle energy possessed by the hot magnetospheric plasma as we know it, is detrimental to the usefulness of frozen-in field lines.

In a cool plasma the particles that are at one time in a common magnetic flux tube will stay together during the subsequent convection. In a hot plasma they do not. This was recognized in the electric field theory of the aurora (Alfvén, 1958) and led to the discovery of the forbidden region phenomenon, which has later been studied by several authors.

For particle energies typical of the hot magnetospheric plasma the differential drift is important. Not only do particles of different species drift differently, but in addition particles of different energy and pitch angle have different drift patterns. Consider for example plasma drifting inward from the plasma sheet. Particles forming an element of plasma back in the plasma sheet will subsequently occupy widely different regions of space (see e.g. Freeman, 1968). Conversely, a volume element of plasma injected from the plasma sheet and observed, say, at the geostationary orbit, will be composed of particles with widely different places of origin.

In principle, one could still use the concept of field line convection, provided the magnetic field lines are electric equipotentials. Then, if the "velocity of magnetic field lines" is defined in the usual way as $\underline{E} \times \underline{B}/B^2$, one would have to keep in mind that no particles (except the rare ones that happen to have zero energy) follow the field line motion. For example, for the drifting plasma clouds observed at the geostationary orbit the $\underline{E} \times \underline{B}/B^2$ drift represents only a small correction to their motion.

3.2 Convection in the presence of nonvanishing $\underline{E}\cdot\underline{B}$

In the low-density magnetospheric plasma, the magnetic field lines may not even be electric equipotentials everywhere. The reason for this will be discussed in §5. If they are not (more precisely if $\underline{B} \times \text{curl } (\underline{B}(\underline{E}\cdot\underline{B})/B^2) \neq 0$, Newcomb, 1958), the concept of moving field lines becomes fundamentally useless, because different parts of the same field line would have to "move" in such a way that they would not remain connected.

For many years now, indications have been accumulating that magnetic-field-aligned electric fields do exist, and recently very direct manifestations of them have been observed (Haerendel *et al.*, 1976). Even more recently, direct measurements at high altitude (up to 8000 km) have demonstrated that at auroral latitudes the electric field does not map as if the geomagnetic field lines were equipotentials and that the field can have magnetic-field-aligned components as large as hundreds of mV/m (Mozer *et al.*, 1977). The measurements show not only that $\underline{E}\cdot\underline{B} \neq 0$ but also that actually $\underline{B} \times \text{curl}(\underline{B}(\underline{E}\cdot\underline{B})/B^2 \neq 0$, so that the condition for frozen-in magnetic field lines is in fact violated.

3.3 The need for improved field models

If the concept of "magnetic field line motion" is abandoned, the convection of particles of any species and energy will have to be studied in terms of the actual magnetic and electric fields. Of these two, the magnetic field has been subject to very extensive observations. But still it is insufficiently known, in the sense that the actual shape of the geomagnetic field lines is very uncertain. This is true already in a geomagnetically quiet situation, and much more so in the presence of strong time variations, e.g. during substorms. This lack of knowledge also has repercussions on the problem of determining the electric field, as was pointed out by Block and Carpenter (1974) and recently by Mozer (1976).

In the case of the electric field, the problem of insufficient empirical basis for quantitative models is much worse. By necessity, the early models of the magnetosphere's electric field were very crude indeed. Often the total electric field was assumed to be a simple superposition of a homogeneous dawn-to-dusk electric field and a field corresponding to a rigid corotation, each of which extended both inside and outside the region of closed flow lines. Such a model has some reasonable general features but of course there is no *a priori* physical basis for a superposition of this kind. The actual electric field is the result of a complicated plasmadynamical interaction, where the spatial distribution of electric fields depends not only on Maxwell's equations (which are linear - hence the superposition principle for vacuum fields) but also on the equations of motion for the plasma (which are non-linear). Furthermore, the exchange of momentum between different flux tubes will depend on electric currents which flow between them, not directly, but by way of the ionosphere as Birkeland currents that may not have any linear relation with the electric field. Unfortunately there seems to have arisen a widespread belief that a superposition principle should apply in this context - much like it was once widely believed that the equivalent current systems of magnetic disturbances had something to do with the flow of the real current.

With the availability of extensive electric field measurements in the magnetosphere, more sophisticated theoretical models have been developed (see e.g. Volland 1975b and references therein). However, as the only detailed data available derives from the ionosphere, any modeling of the electric field in the outer magnetosphere must, for lack of information, be based on assuming magnetic field lines to be equipotentials - an assumption which in view of the results of Mozer _et al_. (1977) is untenable. However, if the field-aligned potential drops are only some tens of kV, and the total potential across the magnetosphere of the order of 100 kV, the problem may not be too serious in the context of the _large-scale_ distribution of electric fields.

The various electric field models that have been developed will not be reviewed here. However, it may be mentioned that Roederer and Hones (1974) have developed a computer code which can be used for calculating particle drifts in any given model of the electric and magnetic field. Time variations can be accomodated too, but whenever a time-dependent magnetic field is involved, the calculation of the induced electric field has, unfortunately, to rely on the assumption $\underline{B} \times \mathrm{curl}\, \underline{B}\, (\underline{E} \cdot \underline{B}/B^2) = 0$.

4 Energization of particles by macroscopic electric fields transverse to the magnetic field

Cosmic plasma has a remarkable capability of energizing charged particles. This is perhaps particularly true in magnetospheres such as those of the Earth and of Jupiter. In any of the mechanisms involved, the actual energization has to be achieved by electric fields. Often this occurs by means of wave fields, but macroscopic electric fields play an important role, too.

4.1 Adiabatic acceleration

The efficiency of macroscopic fields to change the energy of a particle, either by betatron or Fermi acceleration, depends critically on the initial energy of the particle.

Low energy particles will very nearly follow equipotential

surfaces of the electric field and gain no appreciable energy.
Medium energy particles may be defined as those for which the voltage-equivalent of the energy is comparable to the voltage differences available in the magnetosphere. They will be substantially influenced by electric field drift as well as by magnetic gradient and curvature drift and can change their energy substantially in a one-step process (in contrast to the multiple step processes discussed below). The voltage differences available are not well known even for the quasi-static potential field, and much less so for the occasional high fields provided by induction, e.g. during the tail collapse in a substorm. If, for example, the observed "injections" of particles at the geostationary orbit are due to actual inward convection from the plasma sheet, the fields have to be very strong in order to deliver particles over the observed wide energy range.

High energy particles will be dominated by their gradient- and curvature drift. This will in general carry them across some fraction of the electric potential difference available in the magnetosphere and change their energy accordingly. However, for a high-energy particle this energy change means only a relatively small modification of its initial energy. Furthermore, if the particles are trapped, the energy change will be periodic and hence even less interesting. There is, however, one important way in which even high-energy particles can extract energy from the macroscopic electric field in a cumulative way and be substantially energized. If, namely, the magnetospheric electric field (even if it remains conservative at each instant) has random variations on a time scale comparable to the particles' azimuthal drift period. Then some particles will gain, and others will lose, energy repeatedly. The overall result will be a stochastic acceleration and associated diffusive transport (see e.g. Fälthammar, 1968 and references therein). (This process has interesting similarities with the so-called stochastic mode of acceleration in a cyclotron.)

Athough in this case it is the transverse electric field alone that provides the energy, the possibility of decoupling from the iorosphere (by unfreezing of fieldlines) can speed up the transport and associated acceleration by allowing large fluctuations in the transverse electric field to exist relatively unimpeded by the ionospheric "load". In the context of laboratory plasma confinement, numerical simulation by Kamimura and Dawson (1975) has shown that the magnetic mirror field greatly facilitates the decoupling and hence the convective transport. Even a mirror ratio as small as two was found to give a more than threefold increase in the diffusion coefficient. In the magnetosphere, very large mirror ratios are available and may have important consequences.

4.2 Adiabatic acceleration combined with non-adiabatic effects

The convective energization may also be aided by non-adiabatic effects such as pitch angle scattering, which do not in themselves contribute to the energization, but that facilitate the cross-field transport necessary to gain energy from the electric potential field. An example of this is the mechanism discussed by Cole (1971).

The energization of magnetospheric particles that takes place in time-dependent electric fields can be described in terms of a combined betatron and Fermi acceleration. In the stochastic mode just discussed, the energy gain is not limited by the total voltage difference available. It is, however, still subject to the inherent limitations of betatron and Fermi acceleration. (For the betatron effect a limit is set by the maximum relative magnetic field strength as the particle moves inward. The Fermi acceleration is limited by a finite relative reduction in bounce-path length.) These limitations can be circumvented by a cyclic process that is essentially of the same nature as the magnetic pumping long ago proposed for cosmic ray acceleration (Alfvén, 1959). In the case of the magnetosphere this can conceivably be achieved by combining the energizing adiabatic inward radial diffusion and an outward diffusion by essentially energy-conserving non-adiabatic scattering, to a so-called bi-modal diffusion. This is of interest in the case of the radiation belts, and was first proposed by Theodoridis et al. (1968).

More recently, a cyclic process of a related kind has been proposed by Nishida (1976) for particles trapped in Jupiter's radiation belt, and seems to offer a plausible explanation of the very high magnetic moments ($\mu \geq 10^4$ eV/gauss) observed (see van Allen, 1976 and references therein). This process involves a recirculation where particles gain energy through inward radial diffusion. These are scattered to higher L-shells without much energy change. Finally they slide out along high latitude magnetic fieldlines, largely conserving the magnetic moment, which still has its high value when the particles reach the magnetic equator and can start the cycle over again.

In the case of the Earth's magnetosphere, a similar recirculation may be important in a modified version (Shawhan, 1976, personal communication).

4.3 Non-adiabatic acceleration in electric induction fields

Very recently, Pellinen et al. (1977) have shown that in the neighbourhood of a magnetic neutral line in the geomagnetic tail, electric induction fields can bring essentially zero-energy particles up to energies in the MeV range. Essentially, the acceleration consists of two phases. In the first phase, the particle, situated in a weak magnetic field is dominated by the electric force and subject to direct, linear acceleration. As its energy increases to 10-100 keV the magnetic field comes into play and gives the particle a complicated meandering motion, during which it experiences a further energy gain from the non-conservative electric field.

5. Electric fields and currents conduction along the geomagnetic field

5.1 Introduction

The current conduction in a collision dominated plasma is characterized by a local balance between electric and frictional forces. I.e. the momentum gained by the charge carriers in any given volume is lost by collisional friction in the same volume element.

In plasmas of high density (in the terminology of Alfvén and Fälthammar, 1963, Ch. 5.1) this holds transverse as well as parallel to the magnetic field and leads to the ordinary Ohm's law. However, some particles, namely the runaways, are exempted from the local balance, and can carry their momentum largely intact out of the region considered. The fraction of the current that is carried by the runaways is not controlled by the local electric field but by the entire circuit. Only to the extent that this fraction is not large, does the local relation between electric field and current hold.

In plasmas of medium density the electric field vector and the corresponding current density vector need no longer be parallel. The two are now connected by the somewhat more complicated relation called the Generalized Ohm's law, where the conductivity is a tensor rather than a scalar. But it is still a local relation, reflecting a local balance of momentum. The transverse conductivity tends to zero for small as well as for large values of the collision frequency. For conduction parallel to the magnetic field the medium-density plasma is of the same nature as that in a high density plasma.

In a low density plasma, current conduction transverse to the magnetic field is qualitatively similar to that in a medium density plasma. However, along the magnetic field something entirely different happens. Charge carriers can now gain momentum in one place, and lose it by magnetic mirroring in an entirely different part of the magnetic flux tube. Provided the mean free path is long (accounting also for scattering by wave-particle interactions) there need therefore not exist any local relation between electric field and current density.

In regions of magnetically trapped particle populations, magnetic-field-aligned electric fields can exist without generating short circuiting currents. This kind of parallel electric field will be briefly discussed in §5.2.

The magnetic-mirror effect in a low density plasma can also have important effects on current conduction, especially at the relatively large current densities that can occur in auroral Birkeland currents. This will be discussed in §5.3.

Instabilities may lead to generation of wave fields large enough to substantially brake the relative motion of ions and electrons. If the exchange of momentum is large enough it may reinstate a local relation between electric field and current density, that can be described in terms of an "anomalous resistivity". The existence of anomalous resistivity is confirmed by laboratory experiments. The concept is frequently invoked in magnetospheric theory. Some comments on its applicability will be given in §5.4.

Wave-particle interaction has also been proposed to generate a collision-less version of the thermoelectric effect, described in §5.5.

Still another possibility is that instabilities lead to the formation of electric space charge layers, where magnetic field aligned potential drops occur over a small distance. This phenomenon is since long known from laboratory experiments to exist in collision-less plasma. Recently it has also been invoked in the context of magnetospheric phenomena, but like the anomalous resistivity it has not been directly proved to occur in space. Some important features that distinguish space charge layers from anomalous resistivity are discussed in §5.6.

5.2 Parallel electric fields due to different pitch-angle anisotropy of trapped electrons and ions

This mechanism was proposed by Alfvén and Fälthammar (1963, ch. 5.1.3) and studied by Persson (1963, 1966) and Whipple (1977). For a description the reader is referred to §2 of the review by Block and Fälthammar (1976) and to references therein. Here will be given only a qualitative discussion and an illustrating example.

The operation of this mechanism depends on the energy - integrated pitch-angle distributions of electrons and ions being different. Observations of particle distributions at the geostationary orbit indicate that such distributions occur commonly, and that large anisotropies can persist for long times (as evidenced by anisotropic particle clouds being observed to occur periodically during their drift around the Earth). The voltage that can be maintained is related to the voltage equivalent of the particle energies involved. For the drifting clouds at the geostationary orbit this is of the order of kilovolts. The absence of cool plasma is essential.

As a highly idealized example of this mechanism we consider the case of delta-function distributions in energy and pitch angle. Thus let the equatorial pitch angle of electrons and ions be α_{eo} and α_{io} respectively. The local pitch angles α_e and α_i along the magnetic flux tube can then be calculated, and also the local value of E and V at each point. Fig.1 illustrates the result for a case where the electrons have very small pitch angle (like the source-cone electrons observed at ATS-6). The magnetic field has, for simplicity, been taken to be a dipole field. Four cases are shown, corresponding to different combinations of electron and ion energy. Also shown are the intrinsic pitch angle variations, i.e. how the pitch angles would vary with magnetic latitude if there were no parallel electric field.

A simple way in which differential pitch-angle distributions and corresponding parallel electric fields can be established is the following. Suppose that in a certain flux tube particles of one kind, say ions, are introduced. To secure quasineutrality it is necessary not only to pull electrons out from the ionosphere, but also to make the final density distributions of ions and electrons the same. This condition necessitates an electric field, and if an equilibrium exists at all, it will in general be one with a non-vanishing $E_{\parallel}$. Just how large $E_{\parallel}$ will be depends on the details of the injection and is in general extremely hard to calculate.

At least two kinds of injections known to occur in the magnetosphere may cause parallel electric fields of this kind. One kind is the "injections" of plasma taking place during substorms. Even if these were not true injections but just energizations of ambient plasma, they should in general generate a non-vanishing $E_{\parallel}$, because the energization process is likely also to disturb the pitch angle distribution. Another kind of injection takes place whenever a flux tube is invaded by drifting particle clouds such as are often observed at the geostationary orbit.

In mirror-confined laboratory plasma the existence of large potential drops and anisotropic particle distributions has been found by Hopfgarten _et al._ (1968). More recently Geller _et al._ (1974) has given convincing experimental proof of the mechanism

Proposed by Alfvén and Fälthammar. Important recent developments of the theory have been made by Whipple (1977).

5.3 Magnetic-mirror effect on Birkeland currents

Another interesting consequence of the magnetic mirror effect is its influence on current conduction along high latitude magnetic field lines. It has its most pronounced effect in the case of upward current, where it is the electrons that are held back by the mirrors (Rassbach, 1973; Knight, 1973; Lemaire and Scherer, 1974; Lennartsson, 1976, 1977).

We will first consider the adiabatic case and afterwards discuss the role of scattering by plasma turbulence.

The charge carriers available for an upward current from the ionosphere to the plasma sheet are downgoing plasma sheet electrons and upgoing ionospheric ions. The latter are unimpeded by the magnetic mirror force and may therefore have an important role to play (Rassbach, 1973). However, the ions from a stationary ionosphere cannot carry a current larger than their random thermal current which is only a few $\mu A\ m^{-2}$ (at ionospheric level).

Let us now consider upward current conduction by magnetosheath electrons. With typical parameter values such as $T = 10^6 K$ and $n = 10^6 m^{-3}$ the random current at the source is only $1/4\ \mu A\ m^{-2}$. If a substantial part of the random current were utilized to form a net field-aligned current, it would, due to the convergence of the flux tubes, have a density of $0.1 - 1\ mAm^{-2}$ at the ionosphere. That would be enough to account for even the largest current densities observed. However, as we shall see, large magnetic-field-aligned potentials are required in order to draw such currents.

Let the magnetic field strengths at the ionosphere and at the source (in the present example the plasmasheath) be denoted B_I and B_2 respectively, and let i_{e2} be the random electron current at the source. As long as the potential difference is small, plasma sheet electrons will appear near the ionosphere only if their velocity at the source is in the loss cone (which is extremely narrow). Such is the case for the fraction B_2/B_I of the population (if isotropic). However, due to the convergence of the flux tubes, the random current density of these particles at the ionosphere is still i_{e2}. At the ionospheric level this is a very small current density. In order to increase it, a voltage that attracts electrons toward the ionosphere is needed. For small voltages, V, the derivative di/dV is characterized by the thermal voltage equivalent kT_{e2}/e of the particles involved, so that $di/dV = i_{e2}/(kT_{e2}/e)$ which can be rewritten

$$\frac{di}{dV} = \frac{1}{\sqrt{2\pi}} \cdot \frac{e^2 n_{e2}}{m_e v_{Te2}} = \frac{1}{\sqrt{2\pi}} \varepsilon_0 \omega_{pe2}/\lambda_{D2} \qquad (1)$$

where v_{Te2}, ω_{pe2} and λ_{D2} are thermal speed, electron plasma frequency and Debye length at the source.

If any appreciable fraction of the random flux at the source is to be utilized as a charge carrier, i.e. if current densities at the ionospheric level are to be of the order of $i_{e2}B_1/B_2$ (i.e. 0.1 - 1 mAm^{-2}) a voltage is required that can overcome the magnetic mirror repulsion on the bulk of the source particles - not only those with nearly vanishing magnetic moment. The potential required for this is obviously of the order of (B_1/B_2) kT_{e2}/e, which in our example ($T_{e2}=10^6$K, $B_1/B_2 \simeq 10^3$) is 100 kV! As pointed out by Rassbach (1973) the voltage $(B_1/B_2)kT_{e2}/e$ becomes substantial even for fairly cool source electrons. And no potential, however large, can raise the current at ionospheric level above the value $(B_1/B_2)i_{e2}$ (of the order of 0.1 - 1 mAm^{-2} in our example) These features are shown in Fig. 2. For typical parameter values of the plasma sheet the initial shape of the curve represents a <u>conductance per unit area</u> (at ionospheric level) of about 3 $\mu Am^{-2}kV^{-1}$.

The full current-voltage characteristic in the case of a Maxwellian source plasma has been derived by Knight (1973) and Lemaire and Scherer (1974). If we neglect the contribution from outgoing ionospheric electrons, which is cut off at a very small voltage drop, we can write the current density at ionospheric level

$$i = en_{e2}\sqrt{\frac{kT_{e2}}{2\pi m_e}} \cdot \frac{B_1}{B_2}\left[1-\left(1-\frac{B_2}{B_1}\right)\exp\left\{-\frac{eV}{kT_{e2}(B_1/B_2-1)}\right\}\right] \qquad (2)$$

A quantitative diagram of i versus V is given in Fig.3. For typical values of the plasma sheet parameters the current densities observed in auroral Birkeland currents may well require voltage drops of the order of several kilovolts. Notice that the region of proportionality between current densities and voltage (i.e.

constant conductance per unit area) is quite wide (about two decades).

The results discussed so far assume that the source plasma is continually replenished with particles having low magnetic moments. Such replenishment can occur by pitch angle scattering or by continually bringing in new source plasma by transverse convection. If it is not, the current will be choked as the source plasma is depleted of small pitch angle particles.

The distribution of the voltage drop along the magnetic flux tube depends on the entire distribution functions of both electrons and ions along the whole flux tube and is very difficult to calculate.

Notice that the mirror effect imposes a finite (and not very large) conductance per unit area of the magnetic flux tubes. A corresponding conductivity in the sense of a local relation between electric current density and electric field, does not exist.

In the example treated above the source plasma was the magnetosheath. Typical parameters and corresponding conductance (per m^2 in the ionosphere) for different source plasmas are given in Table 1.

So far we have discussed only the completely adiabatic situation and disregarded the effects of scattering. In the case quantitatively discussed above, where the source plasma is continually replenished, the downgoing (although not the upgoing) plasma is isotropic and has a filled loss cone. Then mere scattering, without energy degradation, will not increase the current unless the wave distribution is such as to systematically decrease the particles' magnetic moment even in such a distribution. On the other hand, if the replenishment is incomplete so that the current is partly choked by a deficiency of charge carriers with small pitch angles, then pitch-angle scattering will alleviate the choking and tend to restore the conductance calculated above. If the scattering becomes intense enough, it may itself impede the current and establish an anomalous resistivity.

In this context it is important to notice that the mirror effect can operate even when $i_\parallel$ is far below the random thermal current, so that conditions may not be favourable for intense turbulence.

Although the mirror force can support large potentials along a flux tube, it can locally balance only rather moderate electric forces. As a generous upper limit to the electric field strength that it can balance locally we may take

$$E_{max} = (V + kT_{e2}/e)\left|(\text{grad } B)_\parallel\right|/B \tag{3}$$

where V is the potential drop to the source plasma and T_{e2} is the temperature of the latter. With $V + kT_{e2}/e$ of the order of 10 kV, it means that in the altitude range 2000 to 8000 km a high upper limit is 4 - 15 mVm^{-1}. In this altitude range there have recently been direct measurements (Mozer et al., 1977) of parallel fields as large as hundreds of mV/m (§5.7). Thus the mirror force alone cannot support the large field strength observed in that particular case. However, it may facilitate formation of double layers or electrostatic shocks (Kan, 1976, Lennartsson, 1977) and also influence the structure of such layers due to the (remote) mirroring of electrons trapped below the layer (Lennartsson, 1977).

To determine experimentally to what extent the mirror effect does play a role for maintaining parallel electric fields in the magnetosphere, comprehensive measurements of particle and wave spectra are needed together with measurements of the parallel electric field itself.

5.4 Anomalous resistivity

Anomalous resistivity is the mechanism most commonly invoked to support magnetic field aligned electric fields in space. It has been extensively studied theoretically (although by necessity under highly idealizing assumptions) and by means of computer simulation. Most importantly, its existence has been confirmed in laboratory experiments (see e.g. Schrijver 1973a, b and references therein).

The resistivity of a plasma can be written

$$\eta = \frac{m_e}{e^2 n_e} \nu_{eff} = \frac{1}{\varepsilon_o \omega_{pe}} \cdot \frac{\nu_{eff}}{\omega_{pe}} \tag{4}$$

where ω_{pe} is the electron plasma frequency. This holds generally. In the case of anomalous resistivity, the effective collision frequency ν_{eff} is determined by wave-particle interaction. The magnitude of the dimensionless quantity ν_{eff}/ω_{pe} depends on the distributions of particles as well as waves. Its value for saturated current-driven wave turbulence has been determined in three ways:

(1) from theoretical models of an assumed turbulent state
(2) from numerical simulation
(3) from laboratory experiments.

Several values obtained by experiments and simulation have been compiled by Schrijver (1973b) and are given in Table II.

For comparison with observational data that have recently become available (§5.7) we are in particular interested in the relation between the rms value of the turbulent electric field, E_{rms}, and the dc electric field strength, E, that it may support.

First we note that if a turbulent resistivity is to exist at all, the wave field has to drain momentum from the electrons at a sufficient rate to balance the momentum gain, neE, from the dc electric field (except for a small fraction of the electron population that can be allowed to run away).

A generous upper limit to the rate of momentum imparted from the ac field is neE_{rms}. In any realistic situation with distributed particle velocities and wave fields, it must be much less. Thus we can firmly conclude that any wave field responsible for maintaining anomalous resistivity, in the sense of establishing a local relation of the Ohm's law form $i = E/\eta_{eff}$, must have an intensity such that

$$E_{rms} >> E \tag{5}$$

This is a firm but limited result (we do not know how much larger E_{rms} needs to be than E). We will therefore make the following further considerations.

According to computer simulations (Biskamp et al., 1972)

$$\frac{\nu_{eff}}{\omega_{pe}} = \alpha \frac{\varepsilon_o E_{rms}^2}{2n_e kT_e} \tag{6}$$

with $\alpha = 0.2 - 0.3$ in the state of saturated turbulence. The dc electric field drives a current density

$$i = E/\eta_{eff} \tag{7}$$

which means that the electrons (except the runaways) move with a bulk velocity

$$u_e = \frac{E}{en_e\eta_{eff}} \tag{8}$$

If the turbulence is to be maintained, the bulk velocity must remain about equal to the threeshold velocity for instability even as the electrons are heated by the dissipation of energy. Thus

$$u_e = \beta\sqrt{\frac{2kT_e}{m_e}} \tag{9}$$

where $\beta \lesssim 1$ and depends on the particular instability maintaining the turbulence. For the two stream and Buneman instabilities $\beta \simeq 1$ but for the ion-acoustic and electrostatic ion cyclotron instabilities $\beta < 1$ (typically 0.01 to 0.3 depending on the electron to ion temperature ratio).

It then follows from (4), (6), (8), and (9) that

$$\frac{E_{rms}}{E} = R \simeq \frac{1}{\beta}\sqrt{\frac{\omega_{pe}}{\alpha\nu_{eff}}} \tag{10}$$

With $\alpha = 0.2 - 0.3$, $\beta = 1$ and the values of ν_{eff}/ω_{pe} given in table II, R ranges from 12 to 39, in agreement with the general result (5). If $\beta < 1$, R becomes even larger.

Wave fields of the order of 90 mV/m were observed in the magnetosphere by Fredricks et al. (1973). On the other hand, extensive surveys by the Hawkeye and IMP-6 satellites rarely revealed rms fields larger than 10 mV/m (Gurnett and Frank,1976). According to (10) such wave fields could only support fairly weak dc fields of the order of mV/m or less (which could still integrate to substantial potentials if extended over large bright intervals).

Still another way of looking at the dc field strength is the following. Using (4), (8) and (9) we may write the dc electric field strength in the form

$$E = 2^{1/2}\beta \frac{kT_e}{e\lambda_D} \cdot \frac{\nu_{eff}}{\omega_{pe}} \qquad (11)$$

where $\lambda_D = (\varepsilon_0 kT_e/n_e e^2)^{1/2}$ is the Debye length. According to Table II, ν_{eff}/ω_{pe} is typically of the order of one per cent. Thus, the electric field strength is much less, by about a factor 100, than the thermal voltage equivalent divided by the Debye length. We notice this in preparation for the discussion in §5.6.

One important problem encountered in applying anomalous resistivity to space plasma is that of heat balance. Unlike the mirror effect (§5.3) and space charge layers (§5.6) , the anomalous resistivity leads to local deposition of the dissipated power (primarily in the form of heat given to the local plasma). The magnitude of this power per unit volume is, according to (4) and (6)

$$P = \frac{E^2}{\eta_{eff}} = \alpha^{-1} \cdot \frac{E^2}{E^2_{rms}} \cdot 2\, n_e kT_e \omega_{pe} \qquad (12)$$

P can be quite large. For example, with $\beta=1$ and with the values of ν_{eff}/ω_{pe} from Table I

$$P \gtrsim \frac{n_e kT_e}{50\,\tau_{pe}} \qquad (13)$$

where τ_{pe} is the electron plasma period.

If this power were taken up by the local plasma, it would be very rapidly heated, essentially with a time constant less than a hundred electron plasma periods.

5.5 Thermoelectric effect

Another way in which wave-particle interaction may cause potential drops along magnetic field lines has been proposed by Hultqvist (1971, 1972).

In a collision-dominated plasma, the magnetic-field-aligned thermoelectric field can be written

$$E_{\parallel} = - \frac{k}{en_e} T_e^{-\alpha/2} \frac{d}{ds} (n_e T_e^{1+\alpha/2}) \qquad (14)$$

where α is the thermal diffusion coefficient and has the value 1.4 for a fully ionized plasma of singly charged particles. Only if n_e and T_e vary along the field line in a very special way (such that $n_e T_e^{1+\alpha/2}$ is constant) does this field vanish.

In physical terms the origin of this field is as follows. The electron-ion collision cross section decreases rapidly with energy. At a boundary between a hot and a cold plasma, electrons from the hot plasma meet less resistance when penetrating into the cold plasma than do the cold electrons moving into the hot plasma. I.e. the former have a faster diffusion rate. If there is no net current flowing, a potential barrier has to build up so as to restore the balance by restraining the hot electrons. This barrier must be of the order of magnitude of the voltage equivalent of the hot plasma electrons, and so directed that the hot plasma is positive.

According to Hultqvist, essentially the same kind of phenomenon should be possible in a collisionless plasma. This requires that there is sufficient wave-particle interaction to brake the particle motion substantially and that this braking has a strong energy dependence.

As the thermoelectric effect does not depend on a heavy current, it need not necessarily meet with the power dissipation problem discussed above in the context of anomalous resistivity. It would also allow field-aligned potential drops over much wider areas than where the strong Birkeland currents flow. On the other hand it is important to notice that the thermoelectric field given in (14) may very well coexist with a resistive electric field due to a current. In a collisionless plasma, where the existence of the thermoelectric effect depends on a sufficient wave intensity, conditions for it to occur may be most favourable in the presence of a current.

To evaluate whether the thermoelectric effect plays a role in the magnetosphere will be possible only when we have much better knowledge of the height variations of the total particle distribution function (or at least of density and temperature) and, in particular, of the distribution and properties of particle scattering wave fields.

The total potential that might be supported by thermoelectric effect is related to the volt-equivalent of the temperature of the hot plasma and could hence be several kilovolts. For the local strength of the electric field, equation (5) would be valid here, too, and so the (collisionless) thermoelectric effect would have the same difficulty as anomalous resistivity in locally supporting large field strengths.

5.6 Electric double layers

Electric double layers were first observed in laboratory plasmas a long time ago (Langmuir, 1929). They are distinguished by the following properties

a) quasi-neutrality is locally violated (space-charge layer)
b) the integrated positive and negative space charges nearly cancel, so that the electric field outside is much weaker than inside
c) the potential across the layer is about equal to or larger (often much larger) than the voltage equivalent of the random particle energy of the surrounding plasmas.

Plasma instabilities may develop either into a turbulent state where the plasma exhibits anomalous resistivity or into a double-layer situation with space-charge separation and a localized potential drop. At present it is not possible to tell under what conditions one or the other will happen. This casts a certain shadow over theoretical models where it is simply assumed that the final state of an instability is one of turbulence.

The double layer phenomenon is not easily accessible to theoretical analysis, although simple models based on idealized assumptions have been made already by Langmuir (1929) and recently by Knorr and Goertz (1974). Therefore detailed experimental studies have an important role to play. Torvén (1965, 1968) and Babič and Torvén (1974, 1975, 1976) have made systematic studies of the properties of double layers. Interesting recent experiments have been published by Quon and Wong (1976). For a plasma with $n_e = 10^{14}\ m^{-3}$, $T_e = 4 \cdot 10^4\ {}^oK$ and a current density of about 10 A m^{-2}, they found potential drops of about 15 V over 3 cm ($\approx 20\ \lambda_D$), i.e. a local strength of about 500 V/m.

Until recently all experiments have been made in non-magnetized plasmas, which leaves open the very serious question of the role of material walls.

Experiments in magnetized plasmas are now being made by Torvén. It has been found that in magnetized plasmas, too, large voltage drops occur when the current density becomes so large that the electron drift velocity equals the thermal velocity.

The conductance of the plasma is then suddenly greatly reduced, but recovers after a number of ion plasma periods. The voltage drop is not evenly distributed along the column but concentrated to an acceleration region with very strong electric fields, although this region is not as simple as in the case of non-magnetized plasma.

It may be noted that strong ion waves and an enhanced resistivity also occur, as the current is increased, often as a precursor to the formation of the acceleration region.

Still another important source of understanding of this phenomenon is numerical simulation. Such work has been performed by Goertz and Joyce (1975) who confirmed the existence of double layers as the end result of a current-driven instability.

It can be shown that the thickness L and the average electric field $\overline{E}$ of a double layer with voltage drop V can be written

$$L = \frac{1}{\gamma} \sqrt{\frac{eV}{kT_e}}\ \lambda_D \qquad (15)$$

$$\bar{E} = \frac{V}{L} = \frac{\gamma e\ n_e}{\varepsilon_o}\sqrt{\frac{eV}{kT_e}}\ \lambda_D = \gamma\sqrt{\frac{eV}{kT_e}}\ \frac{kT_e}{e\lambda_D} \tag{16}$$

The factor γ depends on the shape of the space charge distribution, which in turn may depend on V or current i, as discussed further by Shawhan _et al_. (1977). Here we will only note that γ is typically in the range 0.01 to 0.1.

One important feature of the double layer as a means to support a voltage drop is the fact that the power released is _not_ deposited locally but is carried away by the accelerated high-energy electrons down to low altitude where the atmosphere is an efficient sink. (Some heating may of course occur on the way down due to instability of the particle beam.) Thus the problem of energy deposition, which is serious in the case of anomalous resistivity, is essentially absent in the case of double layers.

Double layers and turbulent resistivity are extreme opposites in the sense that in the former case essentially all electrons are runaways, whereas in the latter case only a small fraction is (if the justification for the term resistivity is to remain).

For a recent review on space applications of double layers and related phenomena, the reader is referred to Block (1975) and to the recent papers by Kan (1976), Swift _et al_ (1976), Kan and Akasofu (1976), and Shawhan _et al_ (1977).

5.7 Recent observations of high-altitude electric fields

The first direct observations of electric fields at high altitude above auroral regions have recently become available from a three-component double-probe experiment (Mozer *et al.*, 1977). They have shown that electric fields much larger than those measured in previously explored regions exist, both transverse and parallel to the magnetic field.

Concerning the transverse field, erlier experiments with Ba--injections (Jeffries *et al.*, 1975, Wescott *et al.*, 1975, 1976) have shown the existence of very large drift velocities, corresponding to transverse electric field strengths of hundreds of mV/m. These strong fields are now confirmed by direct measurements and it is found that they often occur in pairs. Such a double structure in the transverse field is of course to be expected as a consequence of field-aligned voltage drops in sheet structures. They may simply be the upward extension of equipotentials "pulled down" by the field-aligned potential drops.

For the parallel field, Ba-experiments have shown the existence of field-aligned potential drops (Wescott *et al.*, 1976; Haerendel *et al.*, 1976). For example Haerendel *et al.* (1976) recorded a field-aligned potential drop of no less than 7.4 kV but could not give a precise value of the electric field strength.

The direct probe measurements by Mozer *et al.* (1977) show that the parallel electric field strength can be very large, hundreds of mV/m. The corresponding potential drops are not measured directly but it can be judged from the surrounding transverse fields that the potential drop below the satellite can only be a few kV. Hence the observed extreme field strengths reported persist only over a small vertical extent (below the satellite) of the order of a few tens of km. They may well represent the first case of electric double layers observed in space (Shawhan *et al.*, 1977).

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Table 1

Source plasma	n_{e2} [1)] cm^{-3}	T_{e2} [1)] ^{o}K	i_{e2} μAm^{-2}	di/dV $\mu Am^{-2}(kV)^{-1}$	B_1/B_2	$(B_1/B_2)(kT_{e2}/e)$ kV
Plasma sheet	1	10^6	0.25	3	10^3	90
Magnetosheath (front)	30	$2 \cdot 10^6$	10	60	$3 \cdot 10^3$	500
Solar wind	8	$1.5 \cdot 10^5$	0.8	60	10^4	130

1) Ref.: Siscoe (1973)

Table II

Values of ν_{eff}/ω_{pe} (from Schrijver, 1973, b)

Author	Method		Value of ν_{eff}/ω_{pe}
Kalinin *et al*	Experiment	(linear)	$1.25-2.5\cdot10^{-2}$
Wharton et *al*	"	(linear)	$0.83-1.25\cdot10^{-2}$
Schrijver	"	(linear)	$0.53\cdot10^{-2}$
Zavioski *et al*	"	(toroidal)	$0.42\cdot10^{-2}$
Hamberger *et al*	"	(linear)	$0.33-1.33\cdot10^{-2}$
Biskamp	Computer simulation		$0.33-0.67\cdot10^{-2}$

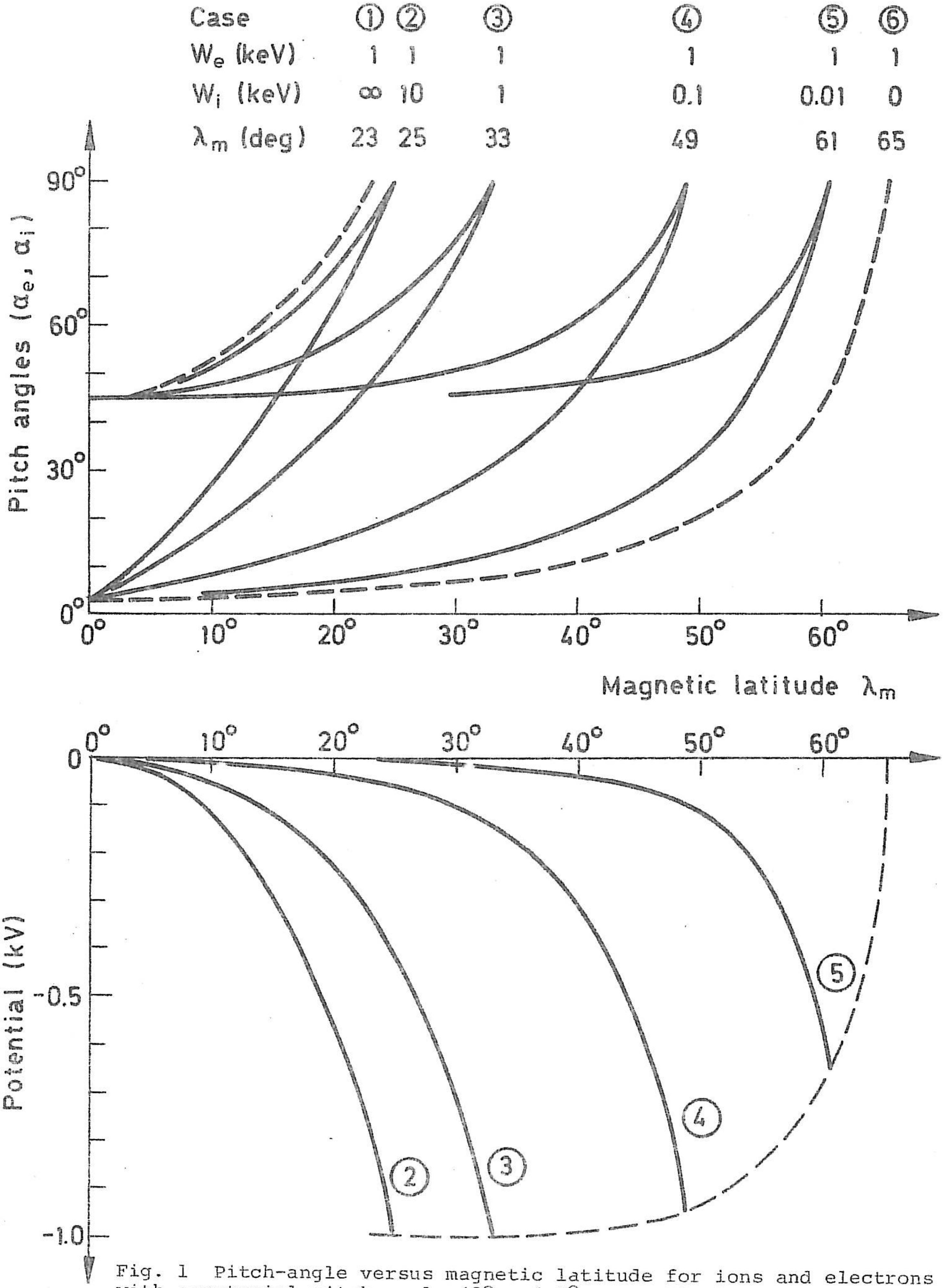

Fig. 1 Pitch-angle versus magnetic latitude for ions and electrons with equatorial pitch-angle 45° and 3° respectively, in a magnetic dipole field. The upper figure shows pitch-angle curves for various combinations of electron and ion energy and the lower figure shows the potentials of the corresponding magnetic-field-aligned electric fields.

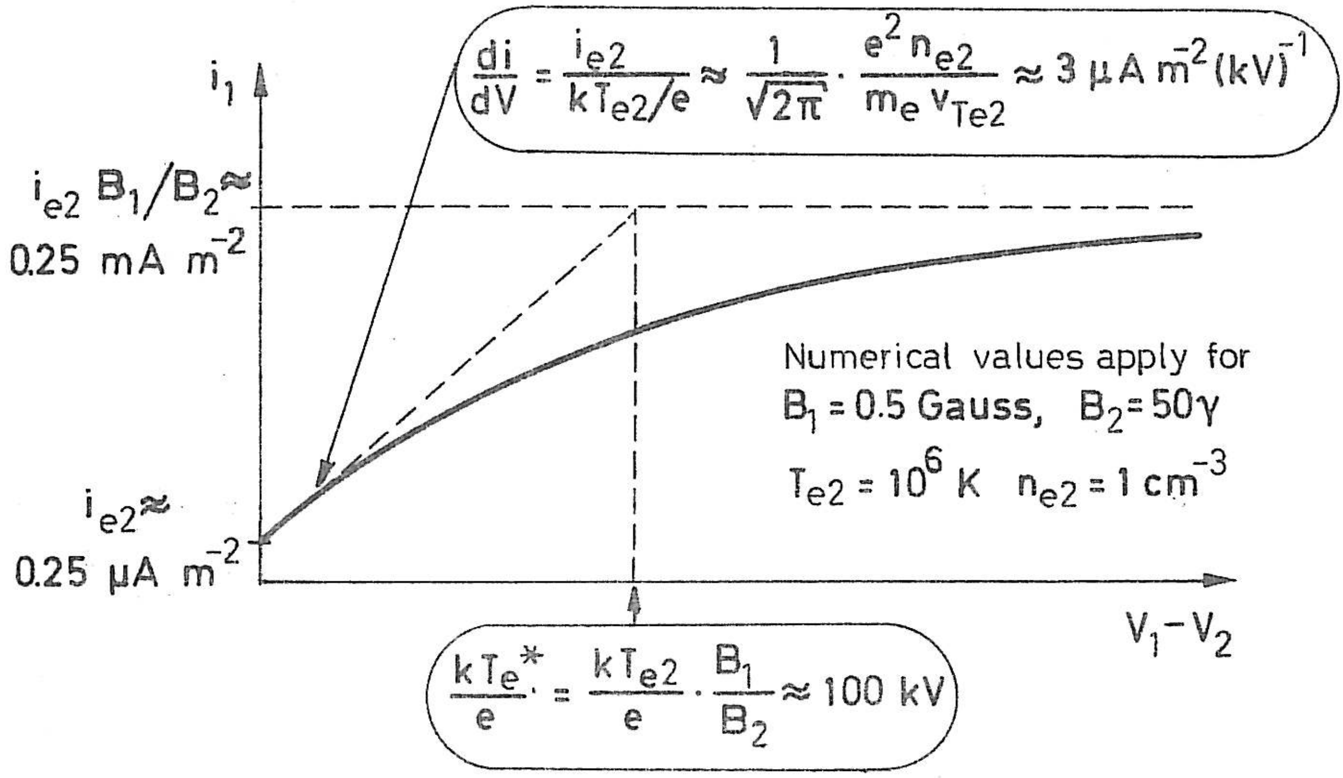

Fig.2 Current-voltage characteristics for upward current carried from the ionosphere (magnetic field B_1) by downgoing electrons from a collisionless source plasma of electron temperature T_{e2} and density n_{e2} in a magnetic field B_2. This characteristic (based on Knight, 1973) applies if the source plasma is continually replenished to remain isotropic and Maxwellian. Otherwise the current will choke as the loss cone is depleted.

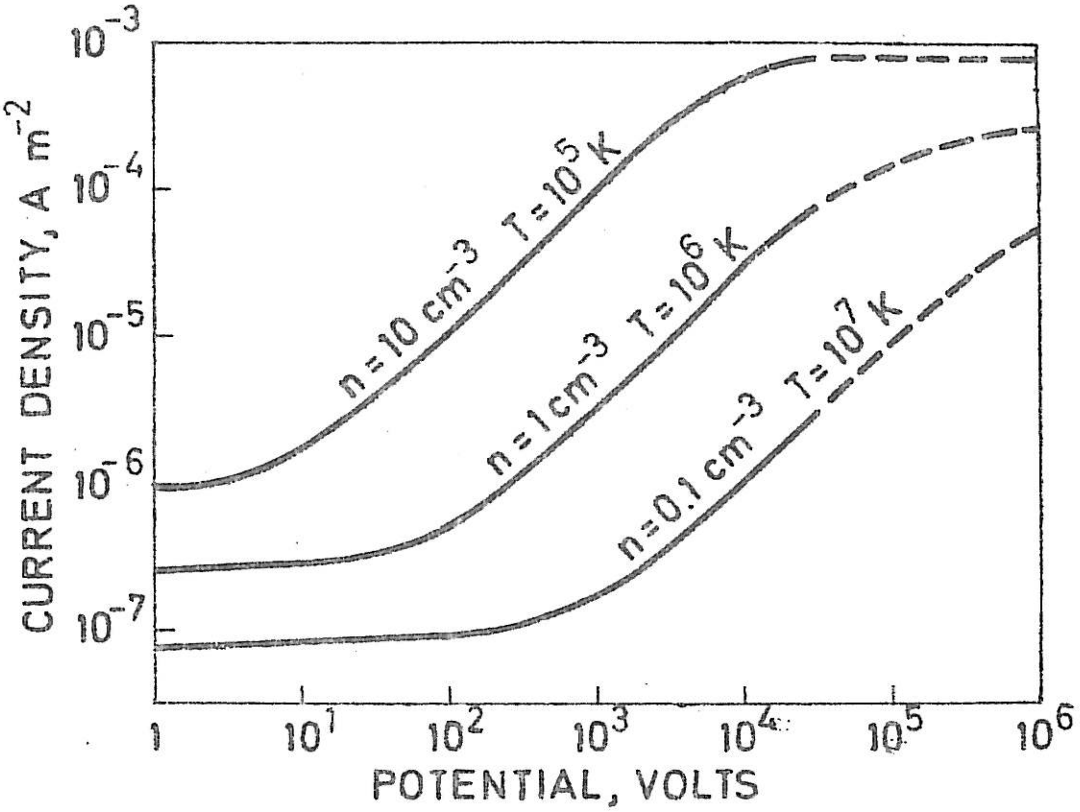

Fig.3 Logarithmic plot of the current-voltage characteristic from Fig.2 for three combinations of source plasma parameters typical for the plasmasheet.

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PROBLEMS RELATED TO MACROSCOPIC ELECTRIC FIELDS IN THE MAGNETOSPHERE

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The macroscopic electric fields in the magnetosphere originate from internal as well as external sources. The fields are intimately coupled with the dynamics of magnetospheric plasma convection. They also depend on the complicated electrical properties of the hot, collisionless plasma. Macroscopic electric fields are responsible for some important kinds of energization of charged particles that take place in the magnetosphere and affect not only particles of auroral energy but also, by multistep processes, trapped high energy particles.

A particularly interesting feature of magnetospheric electric fields is that they can have substantial components along the geomagnetic field, as has recently been confirmed by observation. Several physical mechanisms have been identified by which such electric fields can be supported even when collisions between particles are negligible. Comments are made on the magnetic-mirror effect, anomalous resistivity, collisionless thermoelectric effect and electric double layers, emphasizing key features and differences and their significance in the light of recent observational data
Key words: Electric fields, magnetosphere, anomalous resistivity, magnetic mirror, double layers, space charge layers, thermoelectric effect, convection, electric currents