General Disclaimer

One or more of the Following Statements may affect this Document

- This document has been reproduced from the best copy furnished by the organizational source. It is being released in the interest of making available as much information as possible.

- This document may contain data, which exceeds the sheet parameters. It was furnished in this condition by the organizational source and is the best copy available.

- This document may contain tone-on-tone or color graphs, charts and/or pictures, which have been reproduced in black and white.

- This document is paginated as submitted by the original source.

- Portions of this document are not fully legible due to the historical nature of some of the material. However, it is the best reproduction available from the original submission.

Produced by the NASA Center for Aerospace Information (CASI)

AEROSPACE REPORT NO.
ATR-81(7951)-3

(NASA-CR-168696) KINETIC RESPONSE OF IONOSPHERIC IONS TO ONSET OF AURORAL ELECTRIC FIELDS (Aerospace Corp., El Segundo, Calif.) 38 p HC A03/MF A01 CSCL 04A G3/46
N82-24769
Unclas
21357

Kinetic Response of Ionospheric Ions to Onset of Auroral Electric Fields

Prepared by

Y. T. CHIU
Space Sciences Laboratory
The Aerospace Corporation
El Segundo, Calif. 90245

and

J. R. KAN
Geophysical Institute
University of Alaska
Fairbanks, Alaska 99701

27 July 1981

AUG 7 1981

Prepared for

NATIONAL AERONAUTICS AND SPACE ADMINISTRATION
Headquarters, Washington, D.C. 20545

Contract No. NASW-3434

Laboratory Operations

THE AEROSPACE CORPORATION

MAY 1982
RECEIVED
NASA STI FACILITY
ACCESS DEPT.

Aerospace Report No.
ATR-81(7951)-3

KINETIC RESPONSE OF IONOSPHERIC IONS TO ONSET OF AURORAL ELECTRIC FIELDS

Prepared by

Y. T. Chiu
Space Sciences Laboratory
Laboratory Operations
THE AEROSPACE CORPORATION
El Segundo, Calif. 90245

and

J. R. Kan
Geophysical Institute
University of Alaska
Fairbanks, Alaska 99701

27 July 1981

Prepared for

NATIONAL AERONAUTICS AND SPACE ADMINISTRATION
Headquarters, Washington, D.C. 20545

Contract No. NASW-3434

PRECEDING PAGE BLANK NOT FILMED

Report No.
ATR-81(7951)-3

KINETIC RESPONSE OF IONOSPHERIC IONS
TO ONSET OF AURORAL ELECTRIC FIELDS

Prepared

Y. T. Chiu

J. R. Kan

Approved

J. M. Straus, Head
Atmospheric Sciences Department

G. A. Paulikas, Director
Space Sciences Laboratory

Abstract

By examining the exact analytic solution of a kinetic model of collisional interaction of ionospheric ions with atmospheric neutrals in the Bhatnagar-Gross-Krook approximation, we show that the onset of intense auroral electric fields in the topside ionosphere can produce the following kinetic effects: (1) heat the bulk ionospheric ions to ~ 2 eV, thus driving them up to higher altitudes where they can be subjected to collisionless plasma processes; (2) produce a non-Maxwellian superthermal tail in the distribution function; and (3) cause the ion distribution function to be anisotropic with respect to the magnetic field with the perpendicular average thermal energy exceeding the parallel thermal energy.

PRECEDING PAGE BLANK NOT FILMED

ACKNOWLEDGMENT

The authors benefited from discussions with S. -I. Akasofu, J. M. Cornwall, J. F. Fennell, D. J. Gorney, P. F. Mizera and M. Schulz. We thank R. Lysak for bringing the Russian work to our attention.

CONTENTS

ABSTRACT..... v

ACKNOWLEDGMENT..... vi

I. INTRODUCTION..... 1

II. FORMULATION AND PRELIMINARY CONSIDERATIONS..... 5

III. MODEL OF KINETIC RESPONSE..... 9

IV. SUPERTHERMY AND ANISOTROPY..... 17

V. DISCUSSION AND CONCLUSIONS..... 21

APPENDIX A..... 25

APPENDIX B..... 27

REFERENCES..... 29

I. INTRODUCTION

Ground-based and satellite observations in the latter part of the past decade have brought about a major advance in the understanding of auroral arc formation processes: the electrodynamic interaction between the hot magnetospheric plasma and the cold ionospheric plasma seems to play a central role [e.g., reviews by Akasofu, 1981; Mozer, 1981; Kan and Lee, 1981; Chiu et al., 1981]. Although a major consequence of this electrodynamic interaction is the production of a component of the electric field parallel to the magnetic field for the acceleration of electrons in discrete arc formation, a second, and possibly more far-reaching, consequence of the interaction is the finding that the ionosphere is a significant source of plasmas in the magnetosphere as ionospheric ions are accelerated upwards by auroral electric fields related to substorms [e.g., Shelley et al., 1976; Mizera and Fennell, 1977; Richardson et al., 1981]. This is in addition to the polar wind [Banks and Holzer, 1968; Banks, 1979] which is a significant steady mechanism for transporting ionospheric plasma into the magnetosphere. Actually, the idea that the ionosphere is an active participant in the magnetospheric response to solar-terrestrial activity has had a fairly long history [e.g., Dungey, 1961; Axford and Hines, 1961; Vasyliunas, 1970]. These authors pointed out the importance of the ionosphere in providing the appropriate Pedersen conductivity to close the magnetospheric convection circuit - a role which is crucial in modern theories of magnetosphere-ionosphere coupling [e.g., Chiu and Cornwall, 1980; Kan and Lee, 1980]. In connection with this role of the ionosphere in limiting magnetospheric convection, Joule heating of the ionosphere [Walbridge, 1967; Fedder and Banks, 1972] and the neutral atmosphere [e.g., Ching and Chiu, 1973; Straus and Schulz, 1976] have been considered not only as a thermal

energy source but also as a source to drive ionospheric motions which have an indirect influence upon the magnetospheric convection flow. These studies do not address the question of the ionosphere as a substorm-related source of magnetospheric plasma. In this paper we shall attempt to consider the kinetic properties of the ionospheric plasma as it responds to the onset of an enhanced auroral electric field, and as a source of magnetospheric plasma.

The observational picture of the kinetic properties of such magnetospheric plasmas of ionospheric origin is far from complete; thus, any serious theoretical effort at present must be in the category of "base building." Upward acceleration of ions in auroral electric and magnetic fields is probably not difficult to understand since both the parallel electric field in inverted-V structures and the divergence of the magnetic flux tube favor adiabatic upward acceleration of ionospheric ions such as O^+. Nonadiabatic features of auroral plasmas of ionospheric origin (such as heating and generation of superthermal populations of ion beams and conics) are an entirely different matter. For ionospheric ions, the observations are particularly intriguing since not only are these ions somehow energized to superthermal energies in directions parallel [e.g., Richardson et al., 1981] and perpendicular [e.g., Klumpar, 1979] to the magnetic field, but the processes seem to operate over very wide ranges in energy (~ 6 eV - 10 keV) and in altitude (500 km - 8000 km), and over a wide distribution of local times [Gorney et al., 1981]. Such preliminary observational results clearly indicate the direction of present and future theoretical studies of auroral ionospheric ions: how are superthermal and anisotropic ion populations formed and how are ions energized over some four to five orders of magnitude in energy?

Quite possibly, the answer to these questions may be in-situ wave-particle interactions [e.g., Ungstrup et al., 1979; Okuda and Ashour-Abdalla,

1981]; perhaps future quantitative simulations with more realistic conditions will answer the question of formation of beams (parallel energization) and conics (perpendicular energization) in the observed energies covering the ranges of 6 eV to > 10 keV. Whether such theories are realistic or not is outside the concern and scope of this paper. Rather, we are more interested in the origin of such superthermal populations at the lowest energy ranges (< 6 eV). The question of how special nonadiabatic features are formed out of the cold ionospheric ion population is quite puzzling when one examines the conditions of the ionosphere. First, high-latitude ionospheric temperatures at F-region heights, as measured by the S3-3 satellite, are generally < 2500° K (~ 0.2 eV) [Rich et al., 1979], far less in energy than the superthermal fluxes observed. Second, Kindel and Kennel [1971] concluded that the unstable regions of electrostatic ion cyclotron waves were above the F-maximum and generally in the far topside ionosphere above 1000 km; this makes the relative abundance of O^+ in such nonadiabatic populations (conics and beams) [Ghielmetti et al., 1978] even more puzzling since O^+ must be driven up to such altitudes by some preheating process so that O^+ can at times be the dominant ion at altitudes > 1000 km. Note that in this situation, the abundance of O^+ at altitudes > 1000 km cannot be attributed to escape of H^+ to higher altitudes, as O^+ is also observed to be dominant at the equatorial regions during storm time [Balsiger et al., 1980]. Could it be possible for some kind of pre-heating process to operate in the ionosphere to drive up an intense superthermal (~ several eV) population of ionospheric ions into regions where ion cyclotron waves can act to energize these ions up to tens of keV? Recent observations [Lockwood and Titheridge, 1981] support such hypotheses.

As we have discussed hitherto, the probable occurrence of ion Joule heating in the ionosphere by convection electric fields mapped down to ionospheric heights [Fedder and Banks, 1972] has been well accepted. Since such calculations use a fluid approach, which automatically assumes a Maxwellian form for the distribution function, they cannot tell us about kinetic features of the ion distribution such as pitch angle anisotropy and superthermal populations without going into extremely complex calculations with higher moments [Schunk, 1975]. In view of the necessity to understand how magnetospheric ions may originate from the ionosphere and in view of the predominance of the steady convection electric field as an ionospheric heat source, we are prompted to ask if the auroral electric field may not be the source of preheating which provides the topside with superthermal ions. Since the observed nonadiabatic features are kinetic in character, we are then driven to consider the kinetic response of ionospheric ions to the onset of steady auroral electric fields in a simple collisional kinetic plasma model - the Krook model. Analytically soluble models always oversimplify; thus our main purpose is to gain some insight into kinetic properties of the ion distribution under auroral ionospheric conditions rather than to attempt to "explain" the observed nonadiabatic properties alluded to above. Our work, simple though it may be, differs from previous work on ionospheric heating using a fluid approach [Fedder and Banks, 1972] in that we are required to examine the non-Maxwellian features of the distribution function rather than to assume Maxwellian distributions at all times, as is done in fluid models.

II. FORMULATION AND PRELIMINARY CONSIDERATIONS

Consider a uniform auroral electric field $E_\perp$ imposed upon a horizontally uniform ionosphere consisting of ions, electrons and neutrals whose distribution functions are respectively: f_i, f_e and f_n. In the auroral ionospheric region of interest, where we assume the magnetic field to be vertical and uniform ($\vec{B} = B\hat{z}$), these three species interact through collisions and the total self-consistent electric field $\vec{E}$. The interactions of these species are symbolically formulated in terms of three coupled Boltzmann equations for the distribution function f_k, $k=(i,e,n)$:

$$\hat{L}_k f_k = \sum_{j \neq k} \hat{K}_{kj} (f_k, f_j) \tag{1}$$

where $\hat{L}_k$ is the Boltzmann operator for species k of charge $q_k e$ and mass m_k,

$$\hat{L}_k \equiv \frac{\partial}{\partial t} + \vec{v} \cdot \nabla + (F_k/m_k) \cdot \nabla_v \equiv D/Dt \tag{2}$$

The force F_k consists of a combination of electric, magnetic and gravitational components.

$$\vec{F}_k = m_k \vec{g} + e\, q_k \vec{E} + e\, q_k (\vec{v} \times B\, \hat{z})/c \tag{3}$$

$\hat{K}_{kj}$ is the appropriate binary collision operator. The self-consistent electric field must satisfy Poisson's equation,

$$\nabla \cdot \vec{E} = 4\pi e \int d^3v\, (f_i - f_e) \tag{4}$$

It is well-known that (1) - (4) are very difficult to solve even with the simplest assumptions for $\hat{K}$. In the next section we shall show that a truncated form of these equations with a Krook-type collision term can be solved to give some insight into the kinetic properties of ionospheric ions under the influence of a strong $E_\perp$.

The highly symbolic formulation (1) - (4) nonetheless allows us to discuss certain asymptotic constraints relevant to the problem. If we ignore the fluctuation part of $\vec{E}$ (the auroral DC $E_\perp$ is usually observed to be much larger than the AC electric field), the drivers of non-adiabatic effects are the non-linear collision operators $\hat{K}_{kj}$ which allow energy to be exchanged among the three species. Thus, given sufficient time and no loss of particle species from the system, it is believed that the collisional evolution of the system tends asymptotically toward an isotropic Maxwellian form for the distributions of all the species, although the equilibrium thermal energy of each species need not be the same. This means that the fluid moments approach, with the underlying assumption of isotropic Maxwellian distributions, is valid at times long after onset of $E_\perp$; for ions the time scale is likely to be set by the ion-neutral collision time. In the E-region, where ionic Joule heating effects have been thoroughly studied [Fedder and Banks, 1972] using the fluid equations, the equivalent time scale as determined by the ion-neutral collision time is much less than one second but the neutral response time is ~ 1/2 hour; therefore, we do not expect kinetic features such as superthermy and pitch-angle anisotropy to persist for much longer than the ion-neutral collision time after onset. The situation for the F-region and the topside (> 350 km) is quite different. In this regime, the ion-neutral collision time (> 10 sec) is much longer than the ion gyration period (~ 1/30 sec); therefore, the expected evolution to Maxwellian isotropy takes place

over many gyration periods, permitting kinetic features to persist for tens of seconds or minutes. Can a strong $E_\perp$ drive superthermy and pitch-angle anisotropy in the ionospheric ion population during the evolution period between $E_\perp$ onset and the asymptotic Maxwellian state? It is the intent of this paper to demonstrate that this question can be answered in the affirmative for the Krook model of weakly ionized plasma kinetics.

Before we proceed with the solution of the ionic segment of a Krook-model plasma in uniform electric and magnetic fields, it may be convenient to give an elementary discussion of the origin of the expected kinetic features of superthermy and pitch-angle anisotropy. An imposed $E_\perp$ causes ions to drift and collide with neutral atoms of the thermosphere which act to deflect the uniform drift motion into somewhat random motion. In the initial stages (far from the asymptotic state) the rate of change of ion kinetic energy W_i is not isotropic because the mobilities ($\kappa_\parallel$, $\kappa_\perp$) of ions in the F-region are different for motion parallel and perpendicular to the magnetic field. The perpendicular energy changes according to [Alfvén and Fälthammar, 1963]

$$dW_{i\perp}/dt = e^2 E_\perp^2 \kappa_\perp - \gamma_1 m_i \nu_{in} (W_{i\perp} - W_n)/m_n \tag{5}$$

where $m_i \kappa_\perp = \nu_{in}/(\nu_{in}^2 + \Omega^2)$ [Rishbeth and Garriott, 1969] and $\gamma_1 \approx 2/3$ for the F-region. In (5) m_n and W_n refer to mass and energy of the neutrals; ν_{in} is the ion-neutral collision frequency and Ω is the ion gyrofrequency. The first term on the right hand side of (5) is due to the driving of the electric field and the second term is due to energy loss to the neutrals. The same amount as the energy loss in (5) appears in the corresponding equation for W_n as neutral energy gain. The simultaneous solution of (5) together with the equivalent equation for W_n does not concern us here. The point that the ion population

may be anisotropic soon after $E_\perp$ onset is made by noting that the parallel ion energy is unlikely to increase until sufficient perpendicular drift energy has been scattered through large angles into the parallel direction - probably after many collision times; whereas, $W_{i\perp}$, by virtue of (5), increases in one or a few collision times.

Obviously, the ideas discussed in this section are qualitative. For the rest of this paper, we shall consider the results of a simple solution of an ion component of the Krook model for weakly ionized plasma in steady electric and magnetic fields.

III. MODEL OF KINETIC RESPONSE

The full treatment of the kinetic response of a gravitationally-stratified weakly-ionized and inhomogeneously-magnetized plasma (the F-region ionosphere) to the onset of a strong electric field, as formally presented in the previous Section, is exceedingly difficult to solve. It is the purpose of this Section to initiate a kinetic analysis of this transition region between the collision-dominated E-region and the collisionless magnetosphere with an oversimplified but exactly soluble version of (1) - (4). The simplifications introduced are:

1. The neutral component is assumed to remain in static equilibrium. Thus, the neutrals act as scatterers of ion motion but do not pick up any energy in the process. This is a reasonable approximation if the ionic concentration is sufficiently low or if the average neutral mass is high. Neither condition is strictly valid for the F-region. Fedder and Banks [1972] showed that in a fluid model the motion of the neutrals is an important determinant of the ion heating on time scale greater than about one half hour after electric field onset; since we are dealing with ion response at early times (several collision times after onset), the assumption of immobile neutrals is approximately valid. We hope to relax this restriction in our next stage of kinetic model development.

2. We ignore the collisional influence of electrons upon the dynamic response of ions to steady electric field onset. This is not a bad approximation for early times because ion-electron collisions do not change the ion energy by very much.

3. We ignore the fluctuating (AC) electric field in (4). The AC field driven by auroral electrons may be very important in ion heating, especially in the form of resonance with electrostatic ion cyclotron waves. This effect has been pointed out by others (e.g., Okuda and Ashour-Abdalla, 1981). Here we are not interested in cyclotron heating but, as we have stated, concentrate primarily on how the bulk of auroral ionospheric ions at ~ 0.2 eV may be heated to several eV on the topside. A complete theory of ion conics and beams cannot, however, ignore the effects of the AC electric field.

4. The collision operator for the ion component of (1) is approximated by the simple Krook model [Bhatnagar et al., 1954], which was proposed specifically to study the approach to equilibrium of weakly-ionized collisional plasmas.

$$\hat{K}_{in} = \nu(\vec{x})\,[-f(\vec{x},\vec{v},t) + n(\vec{x},t)\,f_0(\vec{x},\vec{v},t)] \tag{6}$$

where ν is the ion-neutral collision frequency which thermalizes the ion distribution function f towards the assumed isotropic Maxwellian form f_0 [Bhatnagar et al., 1954],

$$f_0 \equiv [m/2\pi\, T(\vec{x},t)]^{3/2} \exp - [m/2T(\vec{x},t)][\vec{v} - \vec{u}(\vec{x},t)]^2 \tag{7}$$

Note that $\int d^3v\, f_0 = 1$. In (6) and (7) n, $\vec{u}$ and T are the density, flow and energy moments:

$$n \equiv \int d^3v\, f \tag{8}$$

$$\vec{u} \equiv \int d^3v\, \vec{v}\, f/n \tag{9}$$

$$T \equiv \int d^3v\, m[\vec{v} - \vec{u}]^2\, f/3n \tag{10}$$

In order for the Krook model to conserve the above three quantities, (6) - (10) have to be self-consistently included in the solution of (1) [Bhatnagar et al., 1954]. The ion-neutral collision frequency ν can depend on n, but we shall assume ν to be a given function of z only. By using the Krook model, it is assumed that ν does not depend on the velocity $\vec{v}$ because the Krook model does not conserve number and energy if $\nu = \nu(\vec{v}, \vec{x})$.

5. The plasma is assumed uniform in the horizontal direction and $\vec{g} = 0$ in (3). Obviously, this simplification ignores auroral spatial scales, but little progress can be made otherwise. The distribution function f can thus depend only on z, the vertical coordinate.

6. The driving DC field $\vec{E}$ is assumed uniform for $t > 0$

$$\vec{E} = \vec{E}_\perp(t) = E_\perp\, \theta(t)\, \hat{x} \tag{11}$$

Since $\nabla \cdot \vec{E} = 0$, (4) implies $\int d^3v\, f = n_0$ where n_0 is the ionospheric electron density - assumed constant. Thus, self-consistency of the electric field (4) and assumption (11) implies $n(\vec{x},t) = n_0$. This constraint can be verified from our solutions.

Our model of kinetic response is defined by (1) - (4) and (6) - (11). With the minor exception of restrictions on $\vec{E}$ and ν, this model is essentially the extension of the Krook model to the case of collisional plasma in uniform electric and magnetic fields. We seek solutions of this model as an initial value problem in which $\vec{E}_\perp(t)$ causes f to evolve with time from an initial isotropic Maxwellian state

$$f(\vec{x}, \vec{v}, t = 0) = n_0 \, [m/2\pi \, T_0]^{3/2} \exp - [m/2T_0] \, v^2 \tag{12}$$

where T_0 is the cold ionospheric temperature ($\sim 2500°$ K): $T(\vec{x},0) = T_0$, $\vec{u}(\vec{x}, 0) = 0$.

Since the assumed force $\vec{F}$ is divergenceless in space, we seek a solution of our model by a phase-space transformation $(\vec{x}, \vec{v}, t) \rightarrow (\vec{x}', \vec{v}', t)$ so that D/Dt in (2) is transformed into $\partial/\partial t$. This is accomplished by the following transformation between the components, labeled (1, 2, 3), of the above vectors:

$$v_3' = v_3 \tag{13}$$

$$\begin{bmatrix} v_1' \\ v_2' \end{bmatrix} = \begin{bmatrix} \cos \Omega t & - \sin \Omega t \\ \sin \Omega t & \cos \Omega t \end{bmatrix} \begin{bmatrix} v_1 \\ v_2 \end{bmatrix} + (cE_\perp/B) \begin{bmatrix} - \sin \Omega t \\ \cos \Omega t - 1 \end{bmatrix} \theta(t) \tag{14}$$

$$x_3' = x_3 - v_3 \, t \tag{15}$$

$$\begin{bmatrix} x_1' \\ x_2' \end{bmatrix} = \begin{bmatrix} x_1 \\ x_2 \end{bmatrix} - \int_0^t d\tau \begin{bmatrix} v_1'(\tau) \\ v_2'(\tau) \end{bmatrix} \theta(\tau) \tag{16}$$

Note that the coordinates $(\vec{x}', \vec{v}')$ are the time-reversed evolution of the coordinates of an ion under $\vec{E}_\perp$ and $B\hat{z}$ starting at $(\vec{x}, \vec{v})$ at $t = 0$; thus

$$\begin{bmatrix} \vec{v}'(\vec{x}, \vec{v}, 0) \\ \vec{x}'(\vec{x}, \vec{v}, 0) \end{bmatrix} = \begin{bmatrix} \vec{v} \\ \vec{x} \end{bmatrix} \tag{17}$$

Further, since the transformation (14) between velocity spaces entails a simple rotation, the volume element is invariant: $d^3v = d^3v'$. The translation in (14) amounts to a shift of origin in velocity space. Under this transformation, the Krook equation [from (1), (2) and (6)]

$$\frac{Df}{Dt} = \hat{K}_{in} \tag{18}$$

becomes (see Appendix A)

$$\partial f(\vec{x}',\vec{v}',t)/\partial t = \nu[z(\vec{x}',\vec{v}',t)]\{n_0 f_0[\vec{x}(\vec{x}',\vec{v}',t),\ \vec{v}(\vec{x}',\vec{v}',t),t] - f(x',\vec{v}',\vec{t})\}. \tag{19}$$

In (19), for the sake of explicitness, we have written $\vec{x}$, $\vec{v}$ as functions of $(\vec{x}', \vec{v}', t)$. Since we shall be working with the $(\vec{x}', \vec{v}', t)$ coordinates we shall hereafter use $\vec{x}$ and $\vec{v}$ to denote functions of $(\vec{x}', \vec{v}', t)$ specified by the inverse transformation of (13) - (16).

The solution to (19), with the initial conditions (12), is

$$f(\vec{x}',\vec{v}',t) = f(\vec{x}',\vec{v}',0)\ e^{-\nu t} + n_0 e^{-\nu t} \int_0^t d\tau \nu e^{\nu t}\ f_0[\vec{x}(\vec{x}',\vec{v}',\tau),\ \vec{v}(\vec{x}',\vec{v}'\tau),\tau] \tag{20}$$

where $\nu \equiv \nu[z(\vec{x}',\vec{v}',t)]$ in general. Note that this integral representation solves (19) generally once $f(\vec{x}',\vec{v}',0)$ and $f_0(\vec{x},\vec{v},\tau)$ are given. General integral representations for $\vec{u}(\vec{x}',t)$ and $T(\vec{x}',t)$ are obtained by substitution of (20) into (9) and (10). Because of assumption 6 (i.e., $\nabla\cdot\vec{E}_\perp = 0$ and (4)), the possible forms chosen for f_0 are constrained by

$$n(\vec{x}',t) = \int d^3v'\ f(\vec{x}',\vec{v}',t) = n_0 \tag{21}$$

We show in Appendix B that our choices (7) and (12) satisfy this constraint (at least for the simple model given below). The integral representation (20) shows explicitly that, as the consequence of collision with neutrals, the initial distribution $f(\vec{x}',\vec{v}',0)$ disappears in a collision time $(1/\nu)$ while a

new convoluted distribution takes its place.

We have, in principle, solved the problem with (20) and (21); however, devising a tractable model so that (20) can be evaluated explicitly is another matter, since the complexity of the problem is now hidden in the transformed f_0 and ν. We have examined a number of models and the most tractable requires that ν = constant, in evaluating the moments $\vec{u}$ and T, (9) and (10), which are crucial in the solution (20) because f_0 depends on $\vec{u}$ and T. For the rest of this paper, we shall restrict our discussion to the "simple" model: ν = constant.

In this "simple" model, the explicit solution is reduced to solutions for $\vec{u}$ and T of the following set of coupled integral equations derived by substitution of (20) into (9) and (10):

$$\vec{u}(\vec{x}'t)\, e^{\nu t} = \int d^3v'\, \vec{v}\, f(\vec{x}',\vec{v}',0)/n_0 + \nu \int_0^t d\tau\, e^{\nu\tau} \int d^3v'\, \vec{v}\, f_0(\vec{x},\vec{v},\tau) \tag{22}$$

$$[3T(\vec{x}',t)/m]\, e^{\nu t} = \int d^3v'\, [\vec{v} - \vec{u}]^2\, f(\vec{x}',\vec{v}',0) + \nu \int_0^t d\tau\, e^{\nu\tau} \int d^3v'\, [\vec{v} - \vec{u}]^2\, f_0\,(\vec{x},\vec{v},\tau) \tag{23}$$

As we have stated earlier, the dependence of $\vec{v}$ and $\vec{x}$ on $(\vec{x}',\vec{v}',t)$ are suppressed in (22) and (23) for brevity. To reduce these integral equations into tractable form we observe that the velocity space integrals can be performed by applying the relations (13) - (16) to (7) and use the identity

$$(\vec{v} - \vec{u})^2 = [v_1' - u_D \sin \Omega t - u_1 \cos \Omega t + u_2 \sin \Omega t]^2$$

$$+ [v_2' + u_D (1 - \cos \Omega t) - u_1 \sin \Omega t - u_2 \cos \Omega t]^2 + [v_3' - u_3]^2 \quad (24)$$

where $u_D = cE_\perp/B$ is the drift speed. Carrying out the velocity space integrals of (22) and (23) is somewhat tedious; for purposes of illustration, we will exhibit only the procedure for (22): $u_3 = 0$ and

$$\begin{bmatrix} u_1 \\ u_2 \end{bmatrix} = (eE_\perp/m) \int_0^t d\tau \, e^{-\nu\tau} \begin{bmatrix} \cos \Omega\tau \\ -\sin \Omega\tau \end{bmatrix}$$
$$= \frac{(eE_\perp/m)}{\nu^2 + \Omega^2} \left\{ \begin{bmatrix} \sin \Omega t & -\cos \Omega t \\ \cos \Omega t & \sin \Omega t \end{bmatrix} \begin{bmatrix} \Omega \\ \nu \end{bmatrix} e^{-\nu t} + \begin{bmatrix} \nu \\ -\Omega \end{bmatrix} \right\} \quad (25)$$

It is interesting to note that (22) is decoupled from (23) because it turns out that $\int d^3v' \, \vec{v} \, f_0 (\vec{x},\vec{v},\tau) = \vec{u} (\vec{x}',\tau)$, giving the simpler results (25). At $t=\infty$, u_1 and u_2 approach the well-known Pedersen and Hall drifts respectively, as we would expect. Using (24) and (25), the integral equation (23) for the thermal energy $3T (\vec{x}',t)/m$ can be written explicitly as

$$(3T/m) = (3 T_0/m) + 2 u_D^2 \Omega^2 \{(1 - e^{-2\nu t})/2 - (1 - e^{-3\nu t})/3\}/(\nu^2 + \Omega^2)$$

$$+ 2u_D^2 \nu \Omega \{-e^{-\nu t} \sin\Omega t + [\nu\Omega(1-e^{-2\nu t} \cos\Omega t) + (2\nu^2 + \Omega^2) e^{-2\nu t} \sin\Omega t]/$$

$$(4\nu^2 + \Omega^2)\}/(\nu^2 + \Omega^2) \quad (26)$$

The thermal energy of the ions is the sum of the initial thermal energy $3T_0/m$, a positive monotonic contribution due to collisional conversion of the electric drift and a term oscillating at the ion cyclotron frequency. The oscillatory term is unimportant in the F-region because it averages to order

$(\nu/\Omega)^2$ which is much less than 1. As $t \to \infty$,

$$(3T/m) \to (3\ T_0/m) + \frac{1}{3}\ u_D^2\ \Omega^2/(\nu^2 + \Omega^2) \tag{27}$$

Thus, about two thirds of the drift kinetic energy $mu_D^2/2$ goes into thermal energy. The exact amount of energy conversion depends on model assumptions; we regard the result of this model only as a guide to ion heating in the F-region. According to (27) then, the thermal energy increase of the bulk of topside F-region ions under the influence of the auroral electrostatic field is roughly

$$(T - T_0) \sim \frac{1}{9}\ mc^2\ (E_\perp/B)^2 \sim 8\ E_\perp^2 \times 10^{-6}\ \text{eV} \tag{28}$$

where $E_\perp$ is in units of mV/m. Thus, if $E_\perp \sim 500$ mV/m [Mozer, 1981], F-region oxygen thermal energy will be ~ 2 eV - sufficient to energize bulk oxygen expansion into the topside auroral ionosphere [Lockwood and Titheridge, 1981], but not sufficient to cause O^+ escape which requires > 10 eV. Note that this mechanism is important only on auroral field lines where $E_\perp$ is large; if we consider the large-scale convection field ($E_\perp < 100$ mV/m) the thermal energy increase is only < 0.1 eV ($< 10^3$ K). This result is in agreement with the fluid calculations of Fedder and Banks [1972]. Further, this mechanism favors thermal energization of oxygen over hydrogen as observed by Ghielmetti et al. [1978].

IV. SUPERTHERMY AND ANISOTROPY

We presented the formulation and solution of the kinetic model in the previous section with a discussion of thermal energy increase of F-region ions. Since T is the thermal energy moment of the entire distribution function, as defined in (10), the above discussion is not applicable to consideration of non-Maxwellian features of the distribution function. These are discussed in this section.

To show that the solution (20) contains a superthermal "tail" which is non-Maxwellian, we need to express $f(\vec{v}',t)$ in an explicit form. This is very difficult, since the time convolution integral of the transformed f_0 is very complicated because $\vec{u}$ and T are complicated functions of time. Instead of giving a detailed numerical study, we shall show that a superthermal non-Maxwellian "tail" appears at early times: $\nu t \ll 1$ and $\Omega t \gg 1$. Since $\Omega t \gg 1$, the oscillatory terms in $\vec{u}$ and T can be ignored because they average to order $(\nu/\Omega)^2$ or smaller. Thus, to order (ν/Ω), (25) and (26) yield

$$u_1 \simeq c(E_\perp/m)(\nu/\Omega) \; ; \; u_2 \simeq - c(E_\perp/m); \; u_3 \simeq 0 \tag{29}$$

$$T/m \simeq T_0/m + \frac{1}{3} u_D^2 \, \nu^2 t^2; \; u_D = cE_\perp/B \tag{30}$$

Defining $T_E/m \equiv \frac{1}{3} u_D^2$, we can write $f_0\,(\vec{v}',t)$ as

$$f_0 \simeq (m/2\pi T_0)^{3/2} \cdot (1 + T_E\, \nu^2 t^2/T_0)^{-3/2} \cdot \exp - \frac{m}{2} (\vec{v}'-\vec{u})^2/(1 + T_E\, \nu^2 t^2/T_0) \tag{31}$$

For $(T_E/T_0)\,(\nu t) \ll 1$, which is somewhat more restrictive than $\nu t \ll 1$, we can expand f_0 above and write an approximate expression of f to order $(\nu t)^3$,

$$f \approx f(t=0)\{e^{-\nu t} + \nu e^{-\nu t} \int_0^t d\tau\, e^{\nu\tau}\{1 + \frac{1}{2}(T_E/T_0^2)(\nu\tau)^2[m(\vec{v}'-\vec{u})^2 - 3T_0]\}$$

$$\approx f(t=0)\{1 + \frac{1}{6}(\nu t)^3\,(T_E/T_0^2)[m(\vec{v}'-\vec{u})^2 - 3T_0] + O[(\nu t)^4]\} \qquad (32)$$

where $f(t=0)$ is the Maxwellian (12). The factor in curly brackets in (32) indicates clearly that f is a non-Maxwellian with a superthermal tail in the superthermal regions of velocity space where $m(\vec{v}'-\vec{u})^2 > 3T_0$. Note that particles in the $m(\vec{v}'-\vec{u})^2 < 3T_0$ regions have been shifted to the superthermal region.

A second important feature of the early-time approximation (32) to the distribution function is that the superthermal factor (curly bracket) is anisotropic because of (29). The anisotropic factor $(\vec{v}'-\vec{u})^2$ is thus roughly

$$(\vec{v}' - \vec{u})^2 \approx [v_1' - c(E_\perp/m)(\nu/\Omega)]^2 + [v_2' + c(E_\perp/m)]^2 + v_3'^2 \qquad (33)$$

Because $\nu/\Omega \ll 1$ in the F-region, we expect the superthermal part of the population in velocity space to be enhanced in the $\vec{E}_\perp \times \vec{B}$ drift direction; i.e., the $\hat{2}$-direction. More discussion on the anisotropy of the superthermal population will be given in the next section.

The property of anisotropy is not limited to the superthermal part of the population. Indeed, we now show that the thermal energy of the bulk of the distribution is also anisotropic with respect to the magnetic field. We split the thermal energy moment, (10) or (23), into perpendicular and parallel components:

$$W_\perp \equiv \int d^3v'\,(\vec{v}_\perp' - \vec{u}_\perp)^2\, f/n_0 \qquad (34)$$

$$W_{\parallel} \equiv \int d^3v' \, (v_3' - u_3)^2 \, f/n_0 \tag{35}$$

$$W \equiv 3T/m = W_{\perp} + W_{\parallel} \tag{36}$$

where W is given explicitly by (26). A tedious calculation gives

$$W_{\parallel} = (2 - e^{\nu t})(T_0/m) + \frac{1}{3} \cdot \frac{u_D^2}{\nu^2+\Omega^2} \{[\Omega^2/3 + 2\nu^2\Omega^2/(4\nu^2+\Omega^2)] -$$

$$[\Omega^2 + 2\nu^2 \, (1-\cos\Omega t)] \, e^{-\nu t} + [\Omega^2 - (\Omega \cos\Omega t + 2\nu \sin\Omega t) \, 2\nu^2\Omega/(4\nu^2+\Omega^2)] \, e^{-2\nu t}$$

$$- (\Omega^2/3) \, e^{-3\nu t}\} \tag{37}$$

From W and $W_{\parallel}$ we can obtain $W_{\perp}$ from (36). The complexity of (26) and (37) due to the rapidly oscillating terms is non-essential. We can set the $\sin\Omega t$ and $(1 - \cos\Omega t)$ terms to zero by averaging over a time long compared to cyclotron time but short compared to collision time and denote the time-averaged thermal moments with an overbar. Further, with the approximations $\nu^2 \ll \Omega^2$ and $3T_0/m \ll u_D^2$, we obtain the thermal anisotropy ratio

$$\bar{W}_{\perp}/(2 \, \bar{W}_{\parallel}) = 1 + \frac{9}{2} \, e^{-\nu t}/(1 - e^{-\nu t}) \; ; \; t \neq 0 \tag{38}$$

where the factor of 2 associated with $\bar{W}_{\parallel}$ compensates for the two degrees of freedom for the perpendicular component. Note that (38) is valid for $t > 1/\Omega$ from onset because we have taken cyclotron averages. For $t=0$, the exact expressions yield $W_{\perp}/(2W_{\parallel}) = 1$, as we would expect. From (38), we note that as $t \to \infty$ the anisotropy again disappears. During the transition phase

($\nu t \sim 1$), the thermal anisotropy can be quite large (~ 3 or ~ 4). In the topside F-region, the collision times can be tens of seconds to minutes, which are of the order of the expected lifetime of auroral acceleration potential structures [Chiu and Schulz, 1978; Chiu and Cornwall, 1980]. Thus, we would expect thermal ions in the topside auroral F-region to show preferential bulk heating in the perpendicular direction.

V. DISCUSSION AND CONCLUSIONS

Because of the necessity to make simplifying assumptions in order to make the model tractable, we do not claim any direct relevance of our model results to observations except in a qualitative sense. However, even at a qualitative level, it is perhaps worthwhile to discuss the implications of our results, which are listed as follows:

1. In the weakly collisional regime of the topside F-region ($\nu \ll \Omega$), a local steady $E_\perp$ (100 - 1000 mV/m) can heat the ions to a temperature of several electron volts, thus driving the bulk of thermal O^+ to high altitudes (~ 1000 km). This Joule heating is not sufficient to drive the bulk O^+ to escape temperature (~ 10 eV).

2. The thermal energy gain scales as $m E_\perp^2$ where m is the mass of the dominant F-region ion. Thus, the mechanism prefers O^+ heating.

3. An initially isotropic Maxwellian distribution is driven to an isotropic asymptotic state with a higher temperature by an external $E_\perp$ imposed on the system. In the transition state, the distribution is non-Maxwellian with a superthermal tail. The ion distribution in this period (in the vicinity of the $\vec{E} \times \vec{B}$ drift energy $\frac{1}{2} m u_D^2$) is anisotropic with respect to the magnetic field.

4. The thermal energy increase of the entire ion population in the transition state is also anisotropic with perpendicular thermal energy averaging several times that of the parallel thermal energy. This anisotropy disappears after several collision times (~ minutes), when the transition from initial to asymptotic states is complete.

Recently Lockwood and Titheridge (1981) discussed, from the standpoint of observed O^+/H^+ transition altitudes, the necessity to assume an ionospheric O^+ heating mechanism to raise the O^+ temperature to several eV in the auroral region. If Mozer's observation [1981] of large $E_\perp$ in the topside ionosphere is confirmed, we believe this can be accomplished by Joule heating in the auroral region. We emphasize that our discussion of Joule heating (heating of the bulk O^+ population) is in qualitative agreement with the calculation of Fedder and Banks (1972) if we apply our results to the large-scale convection electric field which maps to tens of mV/m in the ionosphere, rather than to the auroral electric field. Auroral electric fields of larger magnitude are more localized and do not map through to the lower ionosphere [Rich et al., 1981]. A probable reason for this may be the effect of the ionospheric conductivity profile on the mapping of electric fields inside the ionosphere [e.g. Chiu, 1974]. Joule dissipation of strong electric fields in the topside ionosphere, such as proposed here, may be a second reason for the absence of strong electric fields in the ionosphere.

Optical observations of the auroral 6300 Å line of $O(^1D)$ indicate that the intensity is too high (by about one order of magnitude) to be explained by either electron impact on atomic oxygen and/or by dissociative recombination of O_2^+ [Sharp et al., 1979]. If the bulk of the O^+ ions on the topside can be heated to > 2 eV or if there is sufficient superthermal flux of O^+, increased population of the $O(^1D)$ state can easily be accomplished by the charge exchange interaction between the hot (> 2 eV) O^+ and cold atomic oxygen - the reaction-product atomic oxygen can easily be in the ^1D state [private communication, A. B. Christensen].

Possibly our model for producing superthermal O^+ by auroral electric fields is a first step toward a theory of auroral ion beams and conics. While

we must constantly be reminded of the extreme simplification of the model, the superthermal part of the kinetic-response distribution function of our model possesses essentially the attributes of the conic distribution since the particles that expand upwards along the magnetic field will acquire a parallel velocity from the conservation of the magnetic moment. It remains to be shown, however, that the model superthermal flux at keV energies is sufficient to account for the observed flux. We hope to do a thorough numerical analysis of the distribution in the future. Clearly our model is not intended to deal with high altitude (>> 1000 km) phenomena where, we expect, little collisional influence and wave-particle interaction to be important. Observationally, it would be interesting to see if measurements of low-energy anisotropic superthermal populations of ions at low altitudes (~ 400 km) [Whalen et al., 1978] can be extended to energies as low as several eV, as in Klumpar [1979]. We believe, as do Lockwood and Titheridge [1981], that a complete description of auroral O^+ energization and injection into the magnetosphere must begin with kinetic processes in the F-region ionosphere itself.

After completion of this study, it was brought to our attention that similar ideas and formulation were discussed in a preprint issued by the Space Research Institute of the Soviet Academy of Sciences [Zakharov et al., 1980]. The discussions of the Russian work do not involve solutions of the Krook model as is done here.

PRECEDING PAGE BLANK NOT FILMED

APPENDIX A: TRANSFORMATION OF THE BOLTZMANN OPERATOR

Application of the transformation $(\vec{x},\vec{v},t) \to (\vec{x}',\vec{v}',t)$, (13) - (16), to the components of the Boltzmann operator $\hat{L}_{\{\vec{x},\vec{v},t\}}$, where the curly brackets indicate functionals, yields the sought-for identity

$$\hat{L}_{\{\vec{x},\vec{v},t\}}\, f(\vec{x},\vec{v},t) = \frac{\partial}{\partial t} f(\vec{x}',\vec{v}',t) \tag{A-1}$$

provided the force is given by

$$\vec{F}\,(\vec{v}) = e\,\vec{E}_{\perp} + e(\vec{v} \times B\hat{z})/c \tag{A-2}$$

as assumed in our model.

The calculation is straightforward but somewhat tedious, so we shall only provide the salient points here. Direct differentiation gives

$$\frac{\partial f}{\partial t}\,(\vec{x},\vec{v},t) = \frac{\partial f(\vec{x}',\vec{v}',t)}{\partial t} + \frac{\partial \vec{x}'}{\partial t} \cdot \nabla'\, f(\vec{x}',\vec{v}',t) + \frac{\partial \vec{v}'}{\partial t} \cdot \nabla_{v'}\, f(\vec{x}',\vec{v}',t) \tag{A-3}$$

where ∇' and $\nabla_{v'}$ are gradient operators with respect to $\vec{x}'$ and $\vec{v}'$ respectively. From (13) - (16), we have

$$\frac{\partial \vec{x}'}{\partial t} = -\,\vec{v}(\vec{x}',\vec{v}',t) \tag{A-4}$$

$$\frac{\partial \vec{v}'}{\partial t} = -\,\vec{F}(\vec{v}') \tag{A-5}$$

Next,

$$\vec{v} \cdot \nabla f(\vec{x},\vec{v},t) = \vec{v}(\vec{x}',\vec{v}',t) \cdot \nabla' f(\vec{x}',\vec{v}',t) \tag{A-6}$$

because $\nabla_i = \nabla_i' \cdot \partial x_i/\partial x_i' = \nabla_i'$. Application of (13) - (16) once more gives

$$\vec{F}(\vec{v}) \cdot \nabla_v f(\vec{x},\vec{v},t) = \vec{F}(\vec{v}') \cdot \nabla_{v'} f(\vec{x}',\vec{v}',t) \tag{A-7}$$

The identity (A-1) is obtained by using (A-3) - (A-7).

APPENDIX B: CONSERVATION OF PARTICLES

It is pointed out in Section III that the choice of model initial and asymptotic distribution functions, (12) and (7) respectively, must be self-consistent with the imposed model constraint $\nabla \cdot \vec{E}_\perp = 0$ which amounts to (21) via (4):

$$n(\vec{x}',t) = \int d^3v' \; f(\vec{x}',\vec{v}',t) = n_0 \tag{B-1}$$

where f is given in terms of the forms (7) and (12) by (20). Here we prove (B-1) by direct integration.

<u>Substitution of (20) into (B-1) yields</u> two integrations over d^3v'. The velocity space integration over $f(\vec{x}',\vec{v}',0)$, (12), yields straightforwardly n_0. The velocity space integration over f_0 looks very complicated but inspection of (24) shows that

$$(\vec{v} - \vec{u})^2 = [\vec{v}' - \vec{w}(t)]^2 \tag{B-2}$$

where $\vec{w}(t)$ does not depend on $\vec{v}'$. The Maxwellian form of f_0 now allows $\int d^3v' \; f_0$ to be integrated, yielding a result independent of $\vec{w}$. Thus, (B-1) yields

$$n(x',t) = n_0 \; e^{-\nu t} + n_0 \; \nu \; e^{-\nu t} \int_0^t d\tau \; e^{\nu\tau} = n_0 \tag{B-3}$$

PRECEDING PAGE BLANK NOT FILMED

REFERENCES

Akasofu, S. -I., Morphology of auroral arcs, in Physics of Auroral Arcs, AGU Mono., 1981.

Alfvén, H. and C. G. Fälthammar, Cosmical Electrodynamics, pp. 146-148, Clarendon, Oxford, 1963.

Axford, W. I. and C. O. Hines, A unifying theory of high-latitude geophysical phenomena and geomagnetic storms, Can. J. Phys. 39, 1433, 1961.

Banks, P. M., Magnetosphere-ionosphere interactions, in Solar System Plasma Physics, Vol. II, pp. 57-104, ed. C. F. Kennel, L. J. Lanzerotti and E. N. Parker, North-Holland, New York, 1979.

Banks, P. M. and T. E. Holzer, The polar wind, J. Geophys. Res. 73, 6846, 1968.

Balsiger, H., P. Eberhardt, J. Geiss and D. T. Young, Magnetic storm injection of 0.9 - 16 keV/e solar and terrestrial ions into the high-altitude magnetosphere, J. Geophys. Res. 85, 1645, 1980.

Bhatnagar, P. L., E. P. Gross and M. Krook, A model for collision processes in gases. I: Small amplitude processes in charged and neutral one-component systems, Phys. Rev. 94, 511, 1954.

Ching, B. K. and Y. T. Chiu, Global distribution of thermospheric heat sources: EUV absorption and Joule dissipation, Planet. Space Sci. 21, 1633, 1973.

Chiu, Y. T., Self-consistent electrostatic field mapping in the high latitude ionosphere, J. Geophys. Res. 79, 2790, 1974.

Chiu, Y. T. and J. M. Cornwall, Electrostatic model of a quiet auroral arc, J. Geophys. Res. 85, 543, 1980.

Chiu, Y. T. and M. Schulz, Self-consistent particle and parallel electrostatic field distributions in the magnetospheric-ionospheric auroral region, J. Geophys. Res. 83, 629, 1978.

Chiu, Y. T., J. M. Cornwall and M. Schulz, Effects of auroral-particle anisotropies and mirror forces on high-latitude electric fields, in Physics of Auroral Arcs, AGU Mono., 1981.

Dungey, J. W., Interplanetary magnetic field and the auroral zones, Phys. Rev. Lett. 6, 47, 1961.

Fedder, J. A. and P. M. Banks, Convection electric fields and polar thermospheric winds, J. Geophys. Res. 77, 2328, 1972.

Ghielmetti, A. G., R. G. Johnson, R. D. Sharp and E. G. Shelley, The latitudinal diurnal and altitudinal distributions of upward flowing energetic ions of ionospheric origin, Geophys. Res. Lett. 5, 59, 1978.

Gorney, D. J., A. Clarke, D. Croley, J. Fennell, J. Luhmann and P. Mizera, The distribution of ion beams and conics below 8000 km, J. Geophys. Res. 86, 83, 1981.

Kan, J. R. and L. C. Lee, Theory of imperfect magnetosphere-ionosphere coupling, Geophys. Res. Lett. 7, 633, 1980.

Kan, J. R. and L. C. Lee, Formation of auroral arcs and inverted V precipitations: An overview, in Physics of Auroral Arcs, AGU Mono., 1981.

Kindel, J. M. and C. F. Kennel, Topside current instabilities, J. Geophys. Res., 76, 3055, 1971.

Klumpar, D. M., Transversely accelerated ions: An ionospheric source of hot magnetospheric ions, J. Geophys. Res. 84, 4229, 1979.

Lockwood, M. and J. E. Titheridge, Ionospheric origin of magnetospheric O^+ ions, Geophys. Res. Lett. 8, 381, 1981.

Mizera, P. F. and J. F. Fennell, Signature of electric fields from high and

low altitude particle distributions, Geophys. Res. Lett. 4, 311, 1977.

Mozer, F. S., Electric field and potential structure of auroral arcs, in Physics of Auroral Arcs, AGU Mono, 1981.

Okuda, H. and M. Ashour-Abdalla, Formation of a conical distribution and intense ion heating in the presence of hydrogen cyclotron waves, UCLA preprint, 1981.

Rich, F. J., C. A. Cattell, M. C. Kelley and W. J. Burke, Simultaneous observations of auroral zone electrodynamics by two satellites: Evidence for height variations in the topside ionosphere, J. Geophys. Res., accepted, 1981.

Richardson, J. D., J. F. Fennell and D. R. Croley, Observation of field-aligned ion and electron beams from SCATHA (P78-2), submitted to J. Geophys. Res., 1981.

Rishbeth, H. and O. K. Garriott, Introduction to Ionospheric Physics, p. 131, Academic, New York, 1967.

Schunk, R. W., Transport equations for aeronomy, Planet. Space Sci. 23, 437, 1975.

Sharp, W. E., M. H. Rees and A. I. Stewart, Coordinated rocket and satellite measurements of an auroral event 2: The rocket observations and analysis, J. Geophys. Res. 84, 1977, 1979.

Shelley, E. G., R. D. Sharp and R. G. Johnson, Satellite observation of an ionospheric acceleration mechanism, Geophys. Res. Lett. 3, 645, 1976.

Straus, J. M. and M. Schulz, Magnetospheric convection and upper atmosphere dynamics, J. Geophys. Res. 81, 5822, 1976.

Ungstrup, E., D. M. Klumpar and W. J. Heikkila, Heating of ions to superthermal energies in the topside ionosphere by electrostatic ion cyclotron waves, J. Geophys. Res. 84, 4289, 1979.

Vasyliunas, V. M., Mathematical models of magnetospheric convection and its coupling to the ionosphere, in Particles and Fields in the Magnetosphere, edited by B. M. McCormac, Reidel, Dordrecht, Netherlands, 1970.

Walbridge, E., The limiting of magnetospheric convection by dissipation in the ionosphere, J. Geophys. Res. 72, 5213, 1967.

Whalen, B. A., W. Bernstein and P. W. Daly, Low altitude acceleration of ionospheric ions, Geophys. Res. Lett. 5, 55, 1978.

Zakharov, A. V., V. A. Liperovsky and S. L. Shalimov, On the possible mechanism of ion transverse acceleration in the auroral zone, preprint, Academy of Sciences of the USSR, Space Research Institute, Moscow, 1980.