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The formation of the Venus ionopause: Interaction between the mantle region and the solar wind

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THE FORMATION OF THE VENUS IONOPAUSE: INTERACTION BETWEEN THE MANTLE REGION AND THE SOLAR WIND

by

MARK J. MATNEY

A THESIS SUBMITTED IN PARTIAL FULFILLMENT OF THE REQUIREMENTS FOR THE DEGREE MASTER OF SCIENCE

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ABSTRACT

The solar wind interacts with the non-magnetic planet Venus by processes within the mantle region, located between the upstream shock and the ionosphere. In this region exospheric neutral atoms from Venus interact with electrons and ions in the moving plasma and modify its flow, resulting in a region of sharp ion density gradients at the boundary between the ionosphere and the mantle called the ionopause.

The effect of mass-loading may be simulated by modifying the mass, momentum, and energy conservation equations to include source and loss terms and electromagnetic forces. Using the assumption that the plasma behaves like a fluid, we construct a model to simplify the physics in the mantle. With this model it is possible to generate an oxygen "ledge" in the ion density similar to the observed ionopause. The calculations also show an enhancement in the magnetic field strength above the ionopause as that observed at Venus.
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INTRODUCTION

1.1 HISTORY

Man's study of the planets is perhaps older than history itself. Yet, for the better part of man's history, our knowledge of these distant bodies was limited to ground based observations. The invention of the telescope greatly improved our understanding of the nature of the planets; for only then did we learn that the planets are bodies much like our Earth, ruled by the same physical laws as terrestrial objects. Telescopic observations, however, were but tantalizing glimpses into the wealth of natural phenomena of our Solar System. The planetary sciences took a giant leap forward with the development of the space program and the subsequent ability to make \textit{in situ} observations of various physical properties of the planets. Investigations into the nature of magnetic fields and their associated charged particle plasmas in space opened doors to previously unknown phenomena.

It was discovered that magnetized plasma from the Sun, traveling at supersonic speeds, interacts with the Earth's intrinsic magnetic field, diverting the plasma flow around the planet. This solar wind-magnetospheric interaction has proven to have subtle physical properties that in many cases, despite the simplicity of the basic physical laws at work, often elude our understanding. Some bodies, however, do not have
substantial intrinsic magnetic fields, and the moving plasma appears to interact with the ionized upper atmosphere directly. The planets Venus and Mars, as well as comets and some planetary satellites fall into this category. The planet Venus is the most extensively studied of these bodies, first by flyby Mariner and Venera missions, and now for the last decade by the Pioneer Venus Orbiter. Despite this long history of direct investigation, many of the basic phenomena concerning the solar wind interaction with Venus are poorly understood.

1.2 OVERVIEW OF SOLAR WIND INTERACTION WITH VENUS

Because Venus has an atmosphere and an ionosphere (due to ionization of atmospheric atoms by solar ultraviolet photons) the solar wind does not simply impact the surface as at the Moon, but is mostly diverted around the body of the planet by fluid-like plasma interactions. Using the observational data, many of the basic details of these interactions may be reconstructed (see Figure 1). The supersonic solar wind moving toward Venus first encounters a shock front ahead of the planet. Like a classical shock, the velocity of the supersonic flow drops dramatically through it, raising the density and pressure of the plasma behind it to compensate.

The post-shock region, known as the mantle, resembles that of a conventional hard blunt-body interaction, except that the atmosphere "softens" the interaction (in actuality, the hard planetary surface is the final obstacle, but only by momentum transfer through the thick neutral atmosphere). In many ways, the interaction up to this point resembles that
of the Earth's magnetosphere, where the solar wind is shocked and diverted around the hard obstacle of the magnetic field. At Venus, however, the exosphere has a much stronger effect on the flow that at the Earth. Neutral atoms are ionized in the mantle, changing the mass, momentum, and energy density of the plasma. This effect becomes more pronounced as the flow moves closer to the planet through an exponentially increasing density of atmospheric neutrals, causing the flow parameters to diverge from those characteristic of an adiabatic magnetospheric interaction. Observationally, this region is characterized by low ion densities (sometimes smaller than the detection limits of the ion instruments) and the presence of superthermals - very energetic ions and electrons resulting from very high temperatures or unusual atomic processes. The magnetic fields tied to the flow, after an initial jump through the shock, continue to increase throughout the mantle region, resulting in a phenomenon loosely termed as the Zwan-Wolf effect. By the time the flow reaches the ionosphere, it has slowed to a fraction of its former velocity. The magnetic field increases to provide a considerable portion of the total pressure of the plasma (magnetic plus particle pressures), and in some cases dominates the total pressure.

Observationally, there is a thin transition region between the low-density mantle plasma and the much higher density ionospheric plasma. Whereas the scale heights in the ionosphere proper are several tens of kilometers, typical ion densities in this region change as much as three orders of magnitude through an altitude range of only about 50 kilometers. This ion "ledge" is the ionopause, which appears to mark a steady state transition between one type of plasma flow and another. Some authors
have sought to define the ionopause as a static boundary where the magnetic pressure from above (where it tends to dominate the total pressure) balances the particle pressure that dominates below\(^6\). But, because the name "ionopause" implies a feature in the ion profile, and this "pressure balance ionopause" is often separated from the "density ionopause" by as much as hundreds of kilometers, the density definition is the one adhered to in this work. Work by Cloutier et al.\(^7\) has shown that the whole picture of such a static pressure balance interaction is itself untenable. Instead, the ionosphere maintains a dynamic balance, where the moving plasma apparently transfers momentum through the ionopause and ionosphere to the neutral atmosphere. This work has also been successfully applied to model the solar-wind/Mars interaction\(^8\).

1.3 RESEARCH GOALS

Originally, the present research was intended to create a comprehensive model of the ionopause, detailing its formation, structure, and dynamic behavior. It soon became clear, however, that our knowledge of the processes in the mantle region was inadequate to create such a model for the ionopause. Efforts to fill the gaps in our knowledge resulted in the present model that adds another step to the full understanding of the formation of the ionopause and aids in the formulation of a complete solar wind-Venus interaction model as well as shedding light on similar phenomena at Mars, comets, and certain moons and nebulae where a plasma/neutral-gas interaction dominates.
Simplicity will be the key to this model. Any model that attempts to explain a phenomenon as complex as fluid flow is by necessity going to entail certain simplifying assumptions. Since even the most basic phenomena associated with mass-loaded flows are often not fully understood, this model limits the forms which the plasma parameters may take in order to keep the model as basic as possible (for example, the phase space distribution function of the plasma is assumed to be Maxwellian throughout the mantle region, despite the presence of a significant non-thermal component). Unfortunately, such an approach runs the risk of excluding critical physical properties of the system. The resulting model should, however, lay a solid physics foundation upon which future research may expand and modify the plasma parameters to include such phenomena.

Chapter 2 goes through the derivation of the basic fluid, kinetic, and resulting plasma equations used in calculations of mass-loaded plasma flows. Chapter 3 details the various mass-loading atomic processes important in the Venus mantle and upper ionosphere and how these processes are to be incorporated into the plasma equations. Chapter 4 presents the assumptions used for the various plasma parameters, as well as the boundary conditions and calculation techniques for the model. Chapter 5 gives a comparison of this model to two other magnetospheric models and compares model calculations to actual data from the Pioneer Venus spacecraft.
CHAPTER 2

BASIC EQUATIONS

2.1 THEORY OF FLUID BEHAVIOR

The first step in constructing a model of mass-loaded plasmas is to determine the forms of the conservation equations. These equations resemble the standard equations for hydrodynamics, but are modified to accommodate the addition and subtraction of materials from the flow (cf. equations of Granger, Fluid Mechanics). Because the frequencies of ion-neutral and electron-neutral collisions are much lower than the cyclotron frequencies in the plasma for the temperature and density ranges characteristic of the Venus mantle region, the electromagnetic interactions tend to dominate the dynamics of the charged particles in the plasma. Consequently, these interactions tend to be fluid-like, hence the similarity to hydrodynamics. Collision terms may strongly influence the plasma, but only as effects acting on a material already behaving like a fluid. The neutral exosphere of Venus is therefore effectively decoupled from the plasma, at least in a fluid sense. In essence, the neutral gas is invisible to the plasma fluid. Instead these interactions are manifested by the use of source and loss terms in a fluid that suddenly "sees" new material added to it, seeming to appear from nowhere.

Figure 2 illustrates this mass-loading in more detail. The diagram assumes a plasma moving at velocity \( \mathbf{v} \) in the \( x \) direction with a magnetic field \( \mathbf{B} \) in the \( z \) direction (bold-faced type represents vector quantities) that
is flowing past a stationary neutral atom. If the mean-free-path of the particles is long enough, the neutral atom will be effectively "blind" to the moving plasma. If, however, the atom is suddenly ionized, it will "see" an electric field of \( \mathbf{E} = -v \times \mathbf{B} = -vB \mathbf{y} \). From the frame of the original fixed atom, the ion with charge \( q \) begins to move in the \( \mathbf{E} (y) \) direction, and the path bends toward the \( x \) direction due to the Lorentz force \( \mathbf{F} = q \mathbf{v} \times \mathbf{B} \). The ion cycles around a half circle until it comes to rest two cyclotron radii downstream, at which point the whole cycle is repeated. From the frame of the moving plasma, the ion appears to be added to the flow with velocity \(-v \mathbf{x}\), from which point it executes cyclotron motion, returning to the same relative position each cycle. Such an ion is added with zero energy, yet the plasma "perceives" it as having a certain thermal energy.

A similar situation may occur in a moving classical fluid when a particle at rest is added to it. The new particle is accelerated into the flow by collisions with the other particles, "stealing" a tiny amount of their energy. After the new particle is thermalized, the total amount of energy is the same, but the energy per particle is reduced because the same energy is now shared among more particles.

In the plasma, the fluid behavior is manifested through the electromagnetic fields, thus the ion is accelerated by "stealing" energy from the flow through the fields. The action of these fields causes particles initially at rest to find themselves traveling at new thermal speeds equal to the plasma flow speed once they are added to the flow. Thus, adding an ion to a supersonic flow will, in general, slow the flow and raise the temperature, because the ion's new thermal speed (which is equivalent to
the plasma bulk flow velocity) is higher than that of the average particle in the flow. The flow consequently slows to maintain total energy balance. Similarly, adding an ion to a subsonic flow will cause the temperature to go down, and the flow to speed up, a counter-intuitive result. In general, the ion and its corresponding electron are added with some energy of their own which they carry into the flow. Similarly, particles removed from the flow carry with them some energy and momentum. Correct formulation of the conservation equations should keep proper track of the various effects.

2.2 CONSERVATION EQUATIONS

The continuity equation requires that the time rate of change of the number particles of species \( i \) in a volume \( V \) equals the number flux of particles entering the volume through the surface bounding it plus the net creation rate of particles within \( V \):

\[
\frac{\partial}{\partial t} (\text{number of species } i \text{ particles in volume } V) = (\text{flux of } i \text{ particles in through the surface of } V) + (\text{production rate of } i \text{ particles in } V) - (\text{loss rate } i \text{ particles in } V) \tag{1a}
\]

which becomes

\[
\frac{\partial}{\partial t} \iiint_V n_i \, dV = -\oint_S n_i \mathbf{v}_i \cdot d\mathbf{A} + \iiint_V (p_i - \lambda_i) \, dV \tag{1b}
\]

where \( n_i \) is the number density, \( \mathbf{v}_i \) is the average velocity, and \( p_i \) is the
production rate and $\lambda_i$ the loss rate per unit volume of species $i$. The triple integrals represent the integration over the volume $V$, and the double integrals over the area $S$. Notice that $v_i \cdot dA$ represents the flow rate through this surface element by the species. The first term on the right side may be changed by the Divergence Theorem to read

$$\iint_S n_i v_i \cdot dA = \iiint_V \nabla \cdot (n_i v_i) \, dV. \quad (2)$$

By eliminating the volume integrals, the equation reduces to

$$\frac{\partial}{\partial t} (n_i) = -\nabla \cdot (n_i v_i) + S_i \quad (3)$$

where $S_i = p_i - \lambda_i$, or for the time independent form

$$\nabla \cdot (n_i v_i) = S_i. \quad (4)$$

The mass conservation equation is easily derived from the continuity equation. The bulk velocity of the flow is defined by

$$v = \frac{\sum_i n_i m_i v_i}{\sum_i n_i m_i}. \quad (5)$$

For the present model, I assume the fluid to be well-behaved to the extent that the average velocity of each species is the bulk flow velocity
\[ \mathbf{v}_i = \mathbf{v}_j = \mathbf{v} \quad \text{(6)} \]

such that each continuity equation may be written

\[ \nabla \cdot (n_i \mathbf{v}) = S_i . \quad \text{(7)} \]

Multiplying the continuity equation for each species by its mass and summing over all species

\[ \sum_i \nabla \cdot (m_i n_i \mathbf{v}) = \sum_i m_i S_i \quad \text{(8a)} \]

\[ \nabla \cdot (\sum_i m_i n_i \mathbf{v}) = \sum_i m_i S_i \quad \text{(8b)} \]

which reduces to

\[ \nabla \cdot (\rho \mathbf{v}) = \sum_i m_i S_i . \quad \text{(8c)} \]

Equation (8c), in conjunction with equation (3), represents the general time independent form of the mass conservation equation, where mass is not truly conserved. As long as the production and loss terms are formulated correctly, though, this equation will correctly describe the velocity/density relations of any fluid-like material. The details of the production and loss terms for the conditions at Venus are discussed in chapter 3.

The conservation of momentum is formulated similarly to the continuity equation. The rate of change of the total momentum within a volume is
equal to the flux of momentum through the surface bounding the volume plus the force on the volume by forces on the surface plus the acceleration due to any body forces acting on the fluid in the volume. This is represented by the following equation:

\[
\frac{\partial}{\partial t} (\text{momentum in volume } V) = (\text{flux of momentum in through the surface of } V) + (\text{acceleration by force acting on } S) + (\text{acceleration by force acting on body of } V)
\]  

(9a)

which becomes

\[
\frac{\partial}{\partial t} \iiint_V (\rho \mathbf{v}) \, dV = -\oiint_S (\rho \mathbf{v}) \mathbf{v} \cdot d\mathbf{A} + -\oiint_S \mathbf{P} \mathbf{v} \cdot d\mathbf{A} + \iiint_V \mathbf{F} \, dV
\]

(9b)

where \( \mathbf{F} \) is the body force per unit volume. The pressure term \( \mathbf{P} \) is here shown in the more generalized tensor form, but this model will assume simple isotropic pressure. The surface integrals are reduced into volume integrals by the Divergence Theorem. If the the volume integrals are eliminated, the equation becomes

\[
\frac{\partial}{\partial t} (\rho \mathbf{v}) = - \nabla \cdot (\rho \mathbf{v} \mathbf{v} + \mathbf{P}) + \mathbf{F}.
\]

(10)

We may drop the tensor notation on the pressure, so that the time independent momentum equation is given by
\[ \nabla \cdot (\rho \mathbf{v} \mathbf{v}) + \nabla P = \mathbf{F} . \]  \hspace{1cm} (11)

If the velocity dyad is expanded, this equation becomes

\[ \mathbf{v} (\nabla \cdot \rho \mathbf{v}) + \rho (\mathbf{v} \cdot \nabla) \mathbf{v} + \nabla P = \mathbf{F} . \]  \hspace{1cm} (12)

In standard fluid mechanics texts the first term is eliminated because \( \nabla \cdot \rho \mathbf{v} = 0 \) for classical fluids, which leaves the familiar convective derivative term \((\mathbf{v} \cdot \nabla) \mathbf{v}\).

For the conservation of energy, the rate of change of the energy within a volume is equal to the energy flux into the volume through the sides plus the work done by forces on the surface in addition to energy put into the volume through heat and body forces.

\[
\frac{\partial}{\partial t} (\text{energy in volume } V) = (\text{flux of energy in through the surface of } V) + (\text{work done by forces acting on } S) + (\text{energy put into volume through heat and body forces})
\]

\hspace{1cm} (13a)

which is written
\[
\frac{\partial}{\partial t} \iiint \left( \frac{1}{2} \rho v^2 + U \right) dV = -\iint_{S} \left( \frac{1}{2} \rho v^2 + U \right) v \cdot dA \\
+ \iint_{S} P v \cdot dA + \iiint_{V} Q dV .
\] (13b)

Here \( U \) represents the internal energy per unit volume, and isotropic pressure has been used. \( Q \) is the term for heat addition and work done by body forces. Using the reductions as before, the equation becomes

\[
\frac{\partial}{\partial t} \left( \frac{1}{2} \rho v^2 + U \right) = -\nabla \cdot \left( v \left( \frac{1}{2} \rho v^2 + U + P \right) \right) + Q
\] (14)

or

\[
\nabla \cdot \left( v \left( \frac{1}{2} \rho v^2 + U + P \right) \right) = Q
\] (15)

for the time independent form.

2.3 KINETIC THEORY

The derivations of the various plasma parameters from the kinetic viewpoint warrant closer inspection at this point. For a monatomic or electron species \( i \), the number density is given by

\[
n_i(x) = \iiint d^3u \ f_i(x,u) ,
\] (16a)

the density by
\[ \rho(x) = \sum_i \iiint d^3u \ m_i \ f_i(x,u) = \sum_i m_i \ n_i(x) \]  

(16b)

where \( m_i \) is the species mass, and the species velocity by

\[ v_i(x) = \frac{1}{n_i(x)} \iiint d^3u \ u \ f_i(x,u) \]  

(16c)

where the integrals are over the velocity space \( d^3u \), and \( f_i(x,u) \) is the phase space distribution function for position \( x \) and velocity \( u \). The bulk flow velocity relation from equation (5) may be written

\[ v(x) = \frac{\sum_i \iiint d^3u \ m_i \ u \ f_i(x,u)}{\sum_i \iiint d^3u \ m_i \ f_i(x,u)} = v_i(x) \]  

(17)

if, as a system, all the components move as a single fluid. As for the momentum density,

\[ \sum_i \iiint d^3u \ m_i \ u \ f_i(x,u) = \sum_i m_i \ n_i(x) \ v_i(x) \]

\[ = \sum_i m_i \ n_i(x) \ v(x) \]

\[ = \rho(x) \ v(x) \]  

(18)

The energy density for a moving fluid is somewhat more interesting. At this point it would be instructive to look at a gas with a Maxwellian velocity distribution function, as given by
\[ f_i(x,u) = n_i \left( \frac{m_i}{2\pi k T_i} \right)^{3/2} \exp \left( - \frac{m_i(u(x)-v_i(x))^2}{2kT_i} \right) \]  \hspace{1cm} (19)

where \( k \) is the Boltzmann constant, \( T_i \) is the temperature of the species, and \( v_i \) is the bulk flow velocity. Using this equation, the energy density would then be given by

\[ \sum_i \int \int \int d^3u \left( \frac{1}{2} m_i u^2 \right) f_i(x,u) = \sum_i \frac{1}{2} n_i m_i v_i^2 + \sum_i \frac{3}{2} n_i k T_i. \]  \hspace{1cm} (20)

For an ideal gas with

\[ P_i = n_i k T_i \quad \text{and} \quad P = \sum_i P_i \]  \hspace{1cm} (21)

the energy density is

\[ \sum_i \left( \frac{1}{2} n_i m_i v_i^2 + \frac{3}{2} n_i k T_i \right) = \frac{1}{2} \rho v^2 + \frac{3}{2} P. \]  \hspace{1cm} (22)

The term on the right corresponds to \( P/(\gamma-1) \) for an ideal monatomic gas with \( \gamma = 5/3 \). This is \( U \), the energy per particle per unit volume. Thus, on average, each particle in such a well-behaved gas has an average energy of \( 1/2mv^2 + 3/2kT \).

For a well-behaved gas, the energy conservation equation may be written as
\[ \nabla \cdot \left( \mathbf{v} \left( \frac{1}{2} \rho \mathbf{v}^2 + \frac{P}{\gamma - 1} + P \right) \right) = \nabla \cdot \left( \mathbf{v} \left( \frac{1}{2} \rho \mathbf{v}^2 + \frac{\gamma P}{\gamma - 1} \right) \right) = Q. \] (23)

Notice that the familiar conductive heat flux term has been eliminated owing to the assumption of well-behaved fluid species with similar average velocities.

A few words must be said here about the effects caused by ions with large cyclotron radii. Hot ions, especially oxygen, tend to have cyclotron radii larger than the typical scale length for the fluid interactions. Such situations may cause substantial deviations from fluid-like behavior. This model is intended to reproduce effects observed at Venus, yet still maintain some level of simplicity. In the future it may be advantageous to examine the effect of such non-fluid behaviors, but they are beyond the scope of this work.

### 2.4 ELECTROMAGNETIC FIELDS

In order to form a proper model of the plasma flow at Venus, it is necessary to include the effects of electromagnetic fields on the flow. Most models of Venus have assumed the magnetic field pressures to be small compared to the particle pressures in this region. Spreiter, in his model\(^{10}\), computes the flow of an ideal gas around a blunt body, and then adds in the magnetic field, forcing it to conform to the fluid flow through the frozen-in-flux condition. Such a model does not show the way in which the magnetic field itself may influence the flow. The models of Lees\(^{11}\) and of Zwan and Wolf\(^{12}\) do include the effects of magnetic field pressure, but they only
consider a magnetospheric interaction by an adiabatic plasma without any mass-loading. The modification of magnetic field pressure near the nose of a blunt body at the expense of particle pressure; the so-called Zwan-Wolf effect, has often been included in discussions of the Venus-solar wind interaction; however, no models have attempted to fully examine such an interaction under the conditions at Venus.

The addition of electromagnetic fields to the fluid equations previously derived is relatively straightforward. The equations of continuity are not affected, but the momentum density equation must include the Lorentz force

$$ F \rightarrow F + J \times B $$

(24)

where $J$ is the current density and $B$ the magnetic field. The force due to an electric field is proportional to the net charge of the plasma. In my model, the electric field force term is ignored because the plasma is assumed to maintain charge neutrality. For the energy density equation, Joule heating becomes important, such that

$$ Q \rightarrow Q + J \cdot E $$

(25)

where $E$ is the electric field. The current density $J$ is determined from Ampere's Law:

$$ J = \frac{1}{\mu_0} \nabla \times B, $$

(26)
provided that

$$\nabla \times \mathbf{E} = 0$$

(27)

and

$$\nabla \cdot \mathbf{B} = 0.$$  

(28)

These equations are shown in their time independent forms. The magnetic and electric fields are related by Ohm's law:

$$\mathbf{E} \approx -\mathbf{v} \times \mathbf{B}.$$  

(29)

In actuality, this equation is an over-simplification, because it ignores the effects of collisions and flow geometry (cf. Rossi and Olbert\textsuperscript{13}). Equation (29) is, however, sufficient for the simple model presented here.

2.5 SUMMARY OF EQUATIONS

At this point the equations for a well-behaved plasma fluid may be summarized\textsuperscript{14}:

$$\nabla \cdot (n_i \mathbf{v}) = S_i$$  

(30a)

$$\nabla \cdot (\rho \mathbf{v}) = \sum_i m_i S_i$$  

(30b)
\[ \nabla \cdot (\rho v v) + \nabla P = F + J \times B \]  \hspace{1cm} (30c)

\[ \nabla \cdot (v (\frac{1}{2} \rho v^2 + \frac{\gamma P}{\gamma - 1})) = Q + J \cdot E \]  \hspace{1cm} (30d)

\[ J = \frac{1}{\mu_0} \nabla \times B \]  \hspace{1cm} (30e)

\[ \nabla \times E = 0 \]  \hspace{1cm} (30f)

\[ \nabla \cdot B = 0 \]  \hspace{1cm} (30g)

\[ E = -v \times B . \]  \hspace{1cm} (30h)
CHAPTER 3

MASS-LOADING EFFECTS

3.1 ATOMIC PROCESSES

In order to incorporate the equations for the mass-loaded plasma at Venus, the correct forms of the various source terms must be formulated in detail. The neutral exosphere of Venus consists predominantly of neutral monatomic oxygen and hydrogen$^{15}$, and since the solar wind is composed mostly of electrons and hydrogen ions, the types of interactions possible in this region are somewhat limited.

The first and most familiar reaction that occurs is photoionization, the process by which solar ultraviolet photons (with energy $hv$; $h$ being Planck's constant and $v$ the frequency of the photon) create an ion electron pair from a previously neutral atom, such that

$$
h v + H \rightarrow H^+ + e^- \quad (31a)$$
$$
h v + O \rightarrow O^+ + e^- \quad (hv \geq 13.6 \text{ eV for both reactions}) \quad (31b)
$$

Coincidentally, both hydrogen and oxygen have a first ionization potential of approximately 13.6 eV. The second ionization potential of oxygen is 35.1 eV; consequently, $O^{++}$ photoproduction rates should be far lower than those for $O^+$. Thus, I assume the plasma in the Venus mantle to be composed exclusively of $O^+$, $H^+$, and electrons.
Another chemical process that should occur in the presence of energetic electrons is the impact ionization of neutral species. The reactions important in the mantle region will be

\[ H + e^- \rightarrow H^+ + e^- + e^- \quad (32a) \]
\[ O + e^- \rightarrow O^+ + e^- + e^- . \quad (32b) \]

These energetic electrons come from two sources: thermal electrons in the hot post-shock region and photoelectrons recently created from photoionizations. The photoelectrons will leave the ion with an energy equal to the photon energy minus the ionization potential. If the original photon had an energy greater than twice the ionization potential of the atom, the ejected photoelectron has sufficient energy to ionize a second neutral atom. Because impact ionization cross sections peak for electron energies near 60-100 eV \((T = 10^6 \text{ K})\), this effect should be minor except for the hottest regions of the mantle. Energetic electrons may also excite the neutrals through inelastic collisions, thereby cooling the plasma, but the cross sections for such reactions are much smaller than for impact ionizations, so the effect is small by comparison. Photoionization does not affect the momentum of the plasma because the ions and electrons are added with zero momentum (in the frame of the planet), but the ions and electrons are generally added with some excess energy from the photon. Some of this energy may be lost due to subsequent impact ionization events, as each collision takes an amount of energy from the flow equivalent to the ionization potential of the neutral being ionized. These relations are
computed using the cross sections, solar spectrum (adjusted for Venus), and secondary impact ionization rates from Banks and Kockarts, *Aeronomy*\(^{16}\).

It may be argued that the velocity distribution of newly created ions and electrons will not be Maxwellian. Numerical calculations within the model keep track of the non-thermal deposition rates of the heating terms, but for the flow equations I assume that those particles added to the flow are rapidly thermalized so that the total plasma velocity distribution remains approximately Maxwellian. That is why \(P/(\gamma-1)\) is substituted for the internal energy term \(U\) in the energy conservation equation. This is a weak assumption in the regions of rapid mass-loading because the mass-production rate may exceed the thermalization rate, but it should adequately describe the mass-loaded heating in most of the regions of interest. Any attempt to keep track of the distribution functions of each species greatly increases the complexity of a model, especially as those distributions deviate from a classical Maxwellian. While such calculations may be useful in the future, it is beyond the simplicity desired in this work.

The impact-ionizations from thermal electrons also affect the mass-loading rate and tend to cool the flow. These rates are computed numerically in the model using the relevant cross sections from *Aeronomy*\(^{17}\).

Loss of particles to chemical recombination reactions will be of two forms. The first is three body collisions:
$H^+ + e^- + M \rightarrow H + M \quad (33a)$

$O^+ + e^- + M \rightarrow O + M \quad (33b)$

(where $M$ is some third body), and the second is radiative recombination:

$e^- + H^+ \rightarrow H + \text{hv} \quad (34a)$

$e^- + O^+ \rightarrow O + \text{hv} \quad . \quad (34b)$

Only the radiative recombination reaction rate is large enough in the density and temperature ranges of the Venus mantle region to significantly affect the chemistry therein. This process affects the momentum of the flow by stealing a tiny bit of momentum for each recombination. The recombination force term $F_R$ will look like

$$F_R = -(L_H m_H v) - (L_O m_O v) \quad (35)$$

where each ion of mass $m_i$ lost from the flow takes with it momentum $m_i v$. $L_H$ and $L_O$ are the loss rates and $m_H$ and $m_O$ the masses of hydrogen and oxygen, respectively. Similarly, each recombination takes with it a small amount of energy such that the recombination cooling rate $Q_R$ is

$$Q_R = -L_H \left( \frac{1}{2} m_H v^2 + \frac{kT}{\gamma - 1} \right) - L_O \left( \frac{1}{2} m_O v^2 + \frac{kT}{\gamma - 1} \right) \quad (36)$$

where $T$ is the total temperature of the plasma.$^{18}$
In hot plasmas electrons interacting with ions may cool the plasma by releasing Bremsstrahlung radiation\textsuperscript{19}. Electrons and ions may also collide with neutrals in elastic collisions. Such interactions do not create or eliminate any ion species, but they may be important in momentum and energy transfer between the plasma and the neutral gas. The electron-neutral collision rates and corresponding force and cooling terms are computed numerically in the model\textsuperscript{20}. The ion-neutral elastic collisions tend to be swamped in the mantle by the next process.

The charge-exchange reaction is where an ion passing near a neutral atom may pick up an electron from the neutral and itself become neutralized, while the neutral becomes a new ion to be swept up by the plasma flow. This model uses the four most important of these reactions in the mantle

\begin{align*}
O + O^+ & \rightarrow O^+ + O & (X_{OO}) \\
H + H^+ & \rightarrow H^+ + H & (X_{HH}) \\
O + H^+ & \rightarrow O^+ + H & (X_{OH}) \\
H + O^+ & \rightarrow H^+ + O & (X_{HO}).
\end{align*}

The $X_{ij}$ terms (where ion i is created by neutralization of ion j) will be used to identify each reaction. The cross sections for these interactions are only weakly dependent on interaction velocity, making the reactions strongly temperature dependent. The elastic collisions between ions and neutrals may actually compete with charge-exchange reactions, but for each type of ion-neutral combination, there exists a critical temperature of interaction
below which polarization collisions dominate, and above which charge-exchange reactions dominate. The temperatures encountered in the mantle tend to lie far above this critical temperature for all of the reactions involved, indicating that charge-exchange collisions dominate in this region\textsuperscript{21}. Consequently, ion-neutral elastic collisions are ignored in this model.

Each charge-exchange reaction is like the loss of one ion from the flow and the simultaneous creation of another one. Charge-exchange between two atoms of the same species does not alter the continuity of the flow, but cross charge-exchange between two species does. Momentum is lost from the flow by the loss of the original ion and none is replaced by the new ion, so there is a net drag on the flow. Similarly, energy is lost from the flow by the loss of the original ion and, in general, the added ion contributes less energy than is lost by the first ion. The force and heating terms may be written (approximately) as

\begin{equation}
F_{ij} = -X_{ij} m_j v
\end{equation}

\begin{equation}
Q_{ij} = X_{ij} \left( \frac{kT_{in}}{\gamma - 1} - \left( \frac{1}{2} m_j v^2 + \frac{kT_j}{\gamma - 1} \right) \right).
\end{equation}

$T_{in}$ is the neutral gas temperature of species i, and $T_j$ is the ion temperature of the j ions. In reality the charge-exchange force and heating terms are somewhat more complicated\textsuperscript{22}, but these relations will suffice for this discussion.

In these reactions and in photoionizations as well, it is possible to create oxygen ions with electrons in excited energy levels\textsuperscript{23}. If the plasma is cool
enough, such excited ions will presumably decay either spontaneously or through collisions; but what if the plasma is hot? The added degrees of freedom caused by thermal excitations of these levels changes the value of the adiabatic constant $\gamma$; so that it becomes temperature dependent. The critical ion temperature for having substantial numbers of thermally excited oxygen ions is around $10^4$ K. Fortunately, the ion temperatures tend to fall below this value in the regions where oxygen densities are important, so the excited states may not seriously affect the interaction at Venus.

3.2 MASS-LOADING SOURCE TERMS

The various terms for the plasma fluid equations may now be worked out including the effect from each type of process. In the following equation, $P_i$ refers to the photoionization rate for ion $i$ (including secondary ionizations due to energetic photoelectrons), $X_{ij}$ refers to the charge-exchange rate of the formation of $i$ ions from the neutralization of $j$ ions, $L_i$ will refer to the loss rate of species $i$, and $K_i$ will represent the impact ionization rate of species $i$ from thermal electrons. For oxygen, the production and loss terms are given by

\[
\begin{align*}
    p_0 &= P_0 & \text{(photoproduction rate of oxygen ions)} \\
    &+ X_{\infty} & \text{(self-charge-exchange rate of oxygen)} \\
    &+ X_{OH} & \text{(charge-exchange production rate of oxygen from hydrogen)} \\
    &+ K_O & \text{(impact ionization rate of oxygen atoms)}
\end{align*}
\] (40a)
\[ \lambda_O = L_O \quad \text{(loss rate of oxygen ions)} \]
\[ + X_{OO} \quad \text{(self-charge-exchange rate of oxygen)} \]
\[ + X_{HO} \quad \text{(charge-exchange production rate of hydrogen from oxygen)} \]

\[ S_O = p_O - \lambda_O \]
\[ = P_O - L_O + K_O + X_{OH} - X_{HO}. \]

Note that the \( X_{OO} \) terms vanish for \( S_O \). Similarly, for hydrogen, the production terms are given by;

\[ p_H = P_H \quad \text{(photoproduction rate of hydrogen ions)} \]
\[ + X_{HH} \quad \text{(self-charge-exchange rate of hydrogen)} \]
\[ + X_{HO} \quad \text{(charge-exchange production rate of hydrogen from oxygen)} \]
\[ + K_H \quad \text{(impact ionization rate of hydrogen atoms)} \]

\[ \lambda_H = L_H \quad \text{(loss rate of hydrogen ions)} \]
\[ + X_{HH} \quad \text{(self-charge-exchange rate of hydrogen)} \]
\[ + X_{OH} \quad \text{(charge-exchange production rate of oxygen from hydrogen)} \]
\[ S_H = p_H - \lambda_H \]
\[ = P_H - L_H + K_H + X_{HO} - X_{OH}. \]  
\hspace{1cm} (41c)

As with the oxygen equations, the self charge-exchange terms drop out of the source term \( S_H \).

For the momentum equation, the neutral atmosphere is assumed to be stationary while the plasma has a bulk flow velocity. The force terms would thus be:

\[ F = \rho g \quad \text{(force due to gravity)} \]
\[ -\lambda_H m_H v \quad \text{(loss of momentum from loss of hydrogen)} \]
\[ -\lambda_O m_O v \quad \text{(loss of momentum from loss of oxygen)} \]
\[ -m_e n_e v_{en} v \quad \text{(drag due to electron-neutral collisions)} \]
\hspace{1cm} (42)

where \( g \) is the acceleration due to gravity, \( m_e \) is the electron mass, \( n_e \) the electron density, and \( v_{en} \) the electron-neutral elastic collision rate. The true momentum losses from the different types of collisions are somewhat more complicated; the more accurate forms are used in the model.

The source terms in the energy conservation equation are given by:

\[ Q = \rho g \cdot v \quad \text{(work done by gravity)} \]
\[ -\lambda_H \frac{1}{2} m_H v^2 \quad \text{(loss of flow energy from loss of hydrogen)} \]
\[ -\lambda_O \frac{1}{2} m_O v^2 \quad \text{(loss of flow energy from loss of oxygen)} \]
\[-(X_{HH} + X_{HO} + X_{OH} + X_{OO}) \frac{kT_i}{(\gamma - 1)}\]

(loss of thermal energy from charge-exchange reactions)

\[-(L_H + L_O) \frac{k(T_e + T_i)}{(\gamma - 1)}\]

(loss of thermal energy from recombinations of ions i and electrons e)

\[+ p_H \frac{kT_{H_n}}{(\gamma - 1)}\] (thermal energy from neutral hydrogen H_n)

\[+ p_O \frac{kT_{O_n}}{(\gamma - 1)}\] (thermal energy from neutral oxygen O_n)

\[+ Q_B\] (cooling due to Bremsstrahlung radiation)

\[+ Q_{en}\] (cooling due to electron-neutral collisions)

\[+ Q_K.\] (cooling due to thermal electron impact ionization)

\[(43)\]

As in the momentum force terms, many of the heating terms are much more complex than are shown here. The terms are integrated appropriately in the model.
CHAPTER 4

FORMULATION OF THE MODEL

4.1 FLOW PARAMETERS

There have been several mathematical models devised to recreate the physics of the mantle and upper ionospheric regions of Venus. Most use the Spreiter's model\textsuperscript{24} in which, as mentioned before, he computes the magnetic fields after the rest of the flow is determined. Spreiter adapted techniques originally used for engineering applications in the calculations of shocks around blunt bodies. He is forced by the limitations of his model to assume the plasma to be adiabatic with no mass-loading. Others have used Spreiter's flow profiles \textit{a priori} to investigate the mantle region; Gombosi\textsuperscript{25} has looked at the role of charge-exchange, and Belotserkovskii et al.\textsuperscript{26} have looked at the role of photoionization.

One of the primary motivations of this study was to go "back to the drawing board" and determine what physical processes really do influence the interactions at Venus. The simple calculations of Breus and Krymskii\textsuperscript{27} represent just such a straightforward approach, despite the limitations of their model where they ignored the magnetic fields and assumed the divergence of the flow to be constant in the mantle. Large fluid-dynamics codes, no matter how well they represent the cutting edge of hydrodynamic computational technology, simply give results that are nonsense if the basic physics principles are not addressed and formulated properly.
One of the inspirations for this study comes from the work of Biermann et al.\textsuperscript{28} in which the interaction of a comet with the solar wind was examined by a simple, quasi-one-dimensional cylindrically symmetric model. The flow was computed along the axis of the stagnation flow line; defined as the line of symmetry where the flow moves directly toward the center of the diverting body, which in the Biermann model was the cometary nucleus. Divergence of the flow was calculated by assuming that the flow perpendicular to the stagnation line increased linearly with distance from that line. The pressure gradient perpendicular to the stagnation line was computed by assuming the shape of the isobaric curves, thereby relating the gradient perpendicular to the flow line to the gradient along the flow line. Using these equations, Biermann was able to easily compute the plasma parameters along the stagnation line for a variety of mass-loading and cometary conditions. While there have been numerous other mathematical formulations of comet-solar wind interactions, most lack the elegance and simplicity of this straightforward fluid approach.

It was decided for this model to use spherical coordinates, instead of cylindrical, to take advantage of the tendency of the pressure distribution to conform to a Newtonian $\cos^2 \theta$ profile (where $\theta$ is the colatitudinal angle from the stagnation line). Such a profile is commonly seen between the shock and the surface of a blunt body supersonic interaction (such as a bullet or the nose of a jet)\textsuperscript{29}. Using a solution of this form, as is done in the fluid mechanics papers of Conti\textsuperscript{30} and Kao\textsuperscript{31}, the fluid parameters may be written to lowest order as
\[ v(r,\theta) \cdot \mathbf{r} = v(r) \cos\theta \]  
\[ v(r,\theta) \cdot \Theta = v(r) \sin\theta \]  
\[ P(r,\theta) = P(r) \cos^2\theta \]  
\[ \rho(r,\theta) = \rho(r) \]  
\[ n_i(r,\theta) = n_i(r) \].

Here, \( \mathbf{r} \) and \( \Theta \) are the unit vectors of spherical coordinates (see Figure 3), with the radius \( r \) measured from the center of Venus, and the colatitude \( \theta = 0 \) corresponding to the stagnation flow line. Notice that the velocity equations, to lowest order, correspond to a profile where, in cylindrical \( r, \phi, z \) coordinates, the velocity in the \( z \) direction would be only weakly dependent on the radius near the stagnation line, while the \( \mathbf{r} \) component would vary directly as the radius (cf. Biermann et al.\textsuperscript{32}). The forms of the number and mass density profiles are shown only in the lowest order forms because their colatitudinal derivatives are never required as long as they remain finite. If the system were adiabatic, it might be expected that the mass density would have a profile such that \( \rho = \cos^{\gamma/2} \theta \) in which case the colatitudinal derivative would be zero for \( \theta = 0 \). It would seem reasonable therefore to assume a well-behaved derivative for the more generalized system described here.
4.2 RESULTING FLUID EQUATIONS

Substituting these profiles into the conservation equations (30), and taking the limit $\theta \to 0$ (in order to eliminate most of the higher-order terms), the continuity equation becomes

$$n_i \left( \frac{2}{r} (u + v) + \frac{\partial v}{\partial r} \right) + v \frac{\partial n_i}{\partial r} = S_i \tag{45}$$

where all of the variables are now functions of $r$ only. The mass conservation equation is computed in the same manner, such that

$$\rho \left( \frac{2}{r} (u + v) + \frac{\partial v}{\partial r} \right) + v \frac{\partial \rho}{\partial r} = \sum_i m_i S_i. \tag{46}$$

The $r$-component of the momentum conservation equation may also be easily computed, taking special care with the convective derivative in spherical coordinates, as

$$v \left( \sum_i m_i S_i \right) + \rho v \frac{\partial v}{\partial r} + \frac{\partial P}{\partial r} = F \cdot r. \tag{47}$$

For the $\theta$-component of the momentum equation, substitution yields:

$$\sin\theta \left( u \left( \sum_i m_i S_i \right) + \rho \frac{u}{r} (u + v) + \rho v \frac{\partial u}{\partial r} - \frac{2P}{r} \right) = F \cdot \theta. \tag{48}$$

The entire equation may be divided through by $\sin\theta$, which gives, for $\theta = 0$;
\[ u \left( \sum m_i S_i \right) + \rho \frac{u}{r} (u + v) + \rho v \frac{\partial u}{\partial r} - \frac{2P}{r} = \frac{F \cdot \theta}{\sin \theta}. \]  \hspace{1cm} (49)

For small angles, the \( F \cdot \theta / \sin \theta \) will not be too large, since \( F \propto v \) (see equations (38) and (42)) and, consequently, \( F \cdot \theta \propto u \sin \theta \). The energy conservation equation is also easily computed as:

\[
(\frac{1}{2} \rho \frac{v^2}{\gamma - 1} P) \left( 2 \frac{(u + v)}{r} + \frac{\partial v}{\partial r} \right) \\
+ \left( \frac{1}{2} v^3 \frac{\partial \rho}{\partial r} + \rho v^2 \frac{\partial v}{\partial r} + \frac{\gamma}{\gamma - 1} v \frac{\partial P}{\partial r} \right) = Q. \hspace{1cm} (50)
\]

### 4.3 ELECTROMAGNETIC FIELD PARAMETERS

In order to modify these equations to accommodate plasma properties, the proper form of the magnetic field must be used. In principle, the full solution of the complete equations (30) is straightforward, but in practice, the magnetic field profiles must be decided, at least partly, in advance. Observationally, the magnetic fields in the post-shock region tend to "drape" about the ionosphere of Venus. This effect may be viewed either as the piling up of magnetic fields tied to the stagnating post-shock flow by frozen-in-flux constraints, or as the result of currents induced in the conductive ionosphere by the electric fields associated with the flowing solar wind plasma. In either case, the phenomenon is part of a self consistent system that seeks to divert and/or absorb the mass-loaded solar wind by the bulk of
the planet in the most efficient way possible. The electrodynamical model of Daniell and Cloutier\textsuperscript{33} shows the draped magnetic fields linked to the induced currents running predominantly through the ionosphere in a direction tangential to the planet surface. The flow model proposes that along the stagnation streamline, any component of the magnetic field perpendicular to that line is modified and enhanced by the slowing flow until it swamps any component parallel to the line. Thus the tangential magnetic fields tend to dominate. It is also of interest to note that if the magnetic field lines are parallel to the flow velocity, ions will have difficulty being added to the flow because the accelerating electric field $E = -v \times B$ is zero.

For the purposes of simplicity, this model thus ignores that component of the magnetic field parallel to the stagnation streamline completely, implying a solar wind with magnetic fields perpendicular to its flow direction. Qualitatively, the interaction of such a solar wind would create near-concentric shells of magnetic flux with the lowest shells conforming to the shape of the ionosphere, much like the layers of an onion. The $J \times B$ Lorentz force for such a geometry manifests itself (in the radial direction) in two terms: one for the radial gradients in the magnetic pressure, and one for the magnetic tension due to the curvature of the field lines. Because the flow is already affected by the gradients in the gas pressure, the total "pressure force" is the sum of these three terms:

$$r \cdot \nabla P_p \quad (\text{particle pressure gradients})$$
$$+ r \cdot \nabla P_B \quad (\text{magnetic pressure gradients})$$
$$+ r \cdot T_B \quad (\text{magnetic tension force})$$

(51a)
where

\[ \mathbf{r} \cdot \nabla P \sim \frac{\partial}{\partial r} P \]
\[ \mathbf{r} \cdot \nabla P_B \sim \frac{\partial}{\partial r} B^2 \]
\[ \mathbf{r} \cdot T_B \sim \frac{B^2}{r}. \]

(51b)

In the mantle region, the pressure gradients dominate, while the tension term remains small. If we ignore the tension term the Lorentz force reduces to that of a flat magnetic field profile, which may be visualized as parallel magnetic fields lying in parallel slabs, each perpendicular to the stagnation streamline. This may be considered as a special case of the previous shell model, with the radii of curvature of the shells so great that at any point they appear as planes, much in the same way that the spherical Earth appears flat to an observer on the surface.

Such a magnetic profile greatly simplifies the forms of Maxwell's equations, for

\[ \nabla \cdot \mathbf{B} = 0 \]

(52)

holds everywhere for parallel magnetic field lines. Similarly, the other electromagnetic terms may be easily computed. The radial component of the Lorentz force (24), with the help of Ampere's Law (30e), is given by
\[ \mathbf{r} \cdot (\mathbf{J} \times \mathbf{B}) = -\frac{1}{\mu_0} \left( B \frac{\partial B}{\partial r} + \frac{B^2}{r} \right) \]  

(53)

and the Joule heating term (25) is given by

\[ \mathbf{J} \cdot \mathbf{E} = -\frac{1}{\mu_0} \left( B \frac{\partial B}{\partial r} + \frac{B^2}{r} \right) v. \]  

(54)

The form of Ohm's Law (30h) may also be computed for this geometry with the help of Faraday's Law (30f) to give:

\[ \frac{\partial v}{\partial r} B + v \frac{\partial B}{\partial r} = -\frac{(u + v)}{r} B. \]  

(55)

The angular dependence of the Lorentz force, \( \theta \cdot (\mathbf{J} \times \mathbf{B}) \), presents a further difficulty. The Zwan-Wolf\textsuperscript{34} model uses a gradient of the magnetic pressure in the colatitudinal direction perpendicular to that of the magnetic field direction. This creates an asymmetry in the colatitudinal velocity components. Zwan and Wolf could make such calculations because of the simplifications made possible by the use of adiabatic flow. Such modifications to the model would be constructive in the future; however, the added complication to the Ohm's Law, as well as the addition of another dimension to the velocity profiles, would go beyond the simplicity desired here.

Instead, this model follows more closely that of Lees\textsuperscript{35}. He assumed that the diversion of the flow around the body (in his model, the Earth's magnetosphere) was due to a \( \cos^2\theta \) dependence of the total pressure.
(particle pressure plus magnetic field pressure), instead of the particle pressure only. Such an arrangement would result in a symmetric colatitudinal dependence of the velocities. This implies that the particle pressure gradient parallel to the magnetic field will be steeper than that of a simple hydrodynamic interaction, because the Lorentz force will be strongest in the direction perpendicular to the magnetic field, and much weaker in the parallel direction. Therefore the particle pressure in the direction parallel to the magnetic field will have to "take up the slack." Thus the symmetry of the particle pressure has been sacrificed for symmetry in the velocity profile. Fortunately, the net result of this assumption on the results of the calculations is limited because the region where the flow is most affected by the magnetic fields is at low altitudes, near the ionosphere, whereas a majority of the divergence of the flow occurs at higher altitudes. Still, the effect on the ionopause region of various particle and magnetic field gradients might be of interest in future research.

4.4 FINAL FORM OF PLASMA EQUATIONS

Using the assumption of axisymmetric total pressure gradients, the \( \theta \)-momentum equation (49) may now be written as

\[
\mathbf{u} \left( \sum \mathbf{m}_i \mathbf{S}_i \right) + \rho \frac{\mathbf{u}}{r} (\mathbf{u} + \mathbf{v}) + \rho \mathbf{v} \frac{\partial \mathbf{u}}{\partial r} - 2 \frac{(P + \frac{B^2}{2\mu_0})}{r} = \frac{\mathbf{F} \cdot \theta}{\sin \theta},
\]

(56)

to which may now be added the continuity equations, (45) and (46), to give
\[ n_i \left( \frac{2}{r} (u + v) + \frac{\partial v}{\partial r} \right) + v \frac{\partial n_i}{\partial r} = S_i , \] (57)

\[ \rho \left( \frac{2}{r} (u + v) + \frac{\partial v}{\partial r} \right) + v \frac{\partial \rho}{\partial r} = \sum_i m_i S_i , \] (58)

plus the modified forms of the r-momentum equation (47) and the energy equation (50),

\[ v \left( \sum_i m_i S_i \right) + \rho v \frac{\partial v}{\partial r} + \frac{\partial P}{\partial r} = F \cdot r - \frac{1}{\mu_o} \left( B \frac{\partial B}{\partial r} + \frac{B^2}{r} \right) , \] (59)

\[ \left( \frac{1}{2} \rho v^2 + \frac{\gamma}{\gamma - 1} P \right) \left( 2 \frac{(u + v)}{r} + \frac{\partial v}{\partial r} \right) \]

\[ + \left( \frac{1}{2} v^2 \frac{\partial v}{\partial r} + \rho v^2 \frac{\partial v}{\partial r} + \frac{\gamma}{\gamma - 1} v \frac{\partial P}{\partial r} \right) \]

\[ = Q - \frac{1}{\mu_o} \left( B \frac{\partial B}{\partial r} + \frac{B^2}{r} \right) v , \] (60)

and Ohm's Law (55),

\[ \frac{\partial v}{\partial r} B + v \frac{\partial B}{\partial r} = -\frac{(u + v)}{r} B . \] (61)

The variables \( u, v, P, B, n_i, \) and \( \rho \) are functions of radius only. The source terms \( S_i, F, \) and \( Q \) are functions of these variables in conjunction with the temperatures and densities of the planetary atmosphere at that radius. This
gives six equations and six unknowns: the radial derivatives of $v$, $u$, $P$, $B$, $n_i$, and $\rho$.

4.5 PLASMA AND BOUNDARY CONDITIONS

The upstream solar wind plasma conditions come directly from measurements made by the spacecraft. The OPA (orbital plasma analyzer) experiment gives the solar wind speed and proton density and temperature for each orbit while the orbital magnetometer supplies the solar wind magnetic field strength. As mentioned before, the model assumes that the solar wind enters the shock near the subsolar line with its magnetic fields perpendicular to the flow direction. The pressure of the solar wind is determined by

$$P_{sw} = n_p k (T_e + T_p)$$ (62a)

where the proton density (equal to the electron density) is given by $n_p$, the electron temperature by $T_e$, and the proton temperature by $T_p$. In general, the solar wind is predominantly composed of protons, and the electron temperature tends to be 3-4 times higher than that of the protons. The pressure may thus be given as

$$P_{sw} = n_p k (4.5 T_p).$$ (62b)

The parameters inside the shock are given by the Rankine-Hugoniot magnetohydrodynamic shock jump conditions:
\[ M_s^2 = \frac{\rho_{\text{sw}} v_{\text{sw}}^2}{\gamma P_{\text{sw}}} \]  

(63a)

\[ M_a^2 = \frac{\mu_0 \rho_{\text{sw}} v_{\text{sw}}^2}{B_{\text{sw}}^2} \]  

(63b)

\[ X = \gamma - 1 + \frac{2}{M_s^2} + \frac{\gamma}{M_a^2} \]  

(63c)

\[ \frac{v_0}{v_{\text{sw}}} = \frac{\rho_{\text{sw}}}{\rho_0} = \frac{B_{\text{sw}}}{B_0} = \frac{X + \sqrt{X^2 - 4 \frac{(\gamma - 2)(\gamma + 1)}{M_s^2}}}{2(\gamma + 1)} \]  

(63d)

\[ p_0 = p_{\text{sw}} + \rho_{\text{sw}} v_{\text{sw}} (v_{\text{sw}} - v_0) + \frac{B_{\text{sw}}^2}{\mu_0} (1 - \frac{v_{\text{sw}}}{v_0}) \]  

(63e)

where \( M_a \) and \( M_s \) are the Alfven and sonic mach numbers, respectively. The "sw" subscript refers to solar wind parameters, and the "0" subscript to post-shock parameters. The initial values of \( u \) (recall that \( v \cdot \theta = u \sin \theta \)) depend on the shape of the shock. This model uses a spherical shock, following the development of Biermann et al. It can be shown that just inside such a shock, near \( \theta = 0 \)

\[ u_0 \approx v_{\text{sw}} \frac{R}{R_s} \]  

(64)

where \( R_s \) is the radius of curvature of the shock at radius \( R \) near the
stagnation line. The Spreiter model gives values of $R_s/R \approx 1.2$, which is the value used here.

Once the plasma conditions behind the shock are determined, the model then proceeds to calculate the derivatives at each point and iterate towards the planet using a Runge-Kutta scheme. Basically, the only parameters that need to be adjusted are the size of the steps, the shock radius, the lower boundary radius, and the ratio of ion temperature to electron temperature. Variations from Newtonian pressure contours may also be adjusted to compare their effects on the flow. Model exospheric densities and temperatures for Venus are taken from the work of Nagy and Cravens\textsuperscript{36}.

The actual running of the model is a straightforward shooting-type method. A shock radius is tested, with the result of the velocity profile determining the next shock radius to be tried. If the shock radius is too high, the plasma speed near the ionosphere will go unstable and become exponentially large. Likewise, if the shock altitude is too low, the plasma speed approaches zero too quickly. Thus, by trial and error, the shock altitude is determined that creates a stable profile where the velocity approaches zero asymptotically near the ionosphere. This is the profile that shows realistic plasma parameters all the way through the mantle, and, hopefully, the onset of the ionopause near the end of the calculations. The calculations are concluded at an altitude consistent with the base of the ionopause because the chemical assumptions made for the model break down in the ionosphere proper.
CHAPTER 5

RESULTS AND CONCLUSIONS

5.1 COMPARISONS TO OTHER MODELS

The model as it stands is easily convertible to one that resembles that of Lees\textsuperscript{37}. All mass-loading, force, and heating/cooling terms are temporarily removed from the computer program, and the plasma parameters are adjusted to reflect those of Lees. The calculated example reproduces the overall structure seen in the Lees model. In Figures 4-6, the plasma parameters are shown normalized to their post-shock values. Following the flow from the shock downward; the velocity $v$ of the flow decreases monotonically to zero at the stagnation point. Density in the present model shows a slight increase before decreasing to zero at the stagnation point. This slight increase is not seen in the Lees model, but the loss of particle density is. This region where particle pressure is depleted and magnetic field pressure enhanced is loosely termed the Zwan-Wolf effect. In actuality, Zwan and Wolf modified and extended the calculations of Lees to better explain this region\textsuperscript{38}. The modification of the magnetic field is observed in other models as well.

The Spreiter model\textsuperscript{39} is easily imitated by this model by eliminating not only the presence of the atmosphere, but by also reducing the magnetic field to an insignificant fraction of the total pressure, because Spreiter's model does not include the effect of the magnetic field upon the flow; it is added after the fluid dynamic profiles have already been computed. Figures 7-10
show plasma parameters normalized to their post-shock values. The density and temperatures both increase toward their stagnation values; both are only slightly higher than the Spreiter results. As before, the downward plasma velocity decrease monotonically to zero at the stagnation point. The magnetic field increases to a very large value relative to its post-shock strength near the stagnation point similar to that in the Spreiter model.

5.2 COMPARISONS TO VENUS DATA

Reinserting the mass-loading effects, the model may now be utilized to compare theoretical calculations to plasma data from a particular orbit. Orbit 170 shows structures representative of typical subsolar flow. The OPA instrument gives solar wind conditions for this orbit as $v_{sw} = 400$ km/sec, proton density $n_p = 9$ cm$^{-3}$, and proton temperature $T_p = 4.5 \times 10^5$ K. The OMAG (orbital magnetometer) shows the solar wind magnetic fields to be 12 nT for this orbit. Adding to these parameters a shock radius of $R_s/R = 1.2$ (typical of a Spreiter-type interaction), and setting the ion temperature at 10% of the total plasma temperature, the solar wind/planetary interaction may now be computed. Figure 11 shows the computed velocity profiles, where $v$ (recall $v \cdot r = v(r)\cos\theta$) is the radial velocity downward smoothly approaching zero near the base of the mantle, and the colatitudinal velocity is given by $u$ (recall $v \cdot \theta = u(r)\sin\theta$). In the same region the oxygen density becomes dominant in the plasma, creating an ionopause-like gradient in the total ion density (see Figures 12 and 13). Just above this region, the colatitudinal velocity component representing flow around the planet reaches a maximum, then drops rapidly as the flow enters the oxygen
dominated region. The altitude for this computed oxygen ledge is approximately 75 km lower than the ionopause observed for this orbit by the OIMS (orbital ion mass spectrometer) instrument (Figure 13). The oxygen ledges computed by this model for various solar wind conditions in general do not show the same dynamic response that the observed ionopause exhibits. This may be due to the general breakdown of the fluid assumptions as the plasma nears the ionosphere, or to the rigidity of the axisymmetric assumptions. The Zwan-Wolf model, which uses asymmetric flow velocities, gives higher plasma densities near the stagnation point than does the Lees model. Because the ionopause is a structure observed in the density profile, any changes in the ion densities due to different velocity assumptions could affect the calculations of those densities. Such future work with this model may better recreate the dynamics of the ionopause formation at Venus.

The model does show a plasma density depletion region above the oxygen ledge that does not appear in the observations and is present in the calculations for a wide range of solar wind conditions. It tends to show ion densities less than 10 cm⁻³, below the limit of sensitivity of the ion instruments aboard the Pioneer Venus spacecraft. The plasma depletion is associated with a corresponding increase in magnetic field pressure (Figure 14). The computed magnetic field reaches a maximum of 74 nT; somewhat less than the observed 106 nT maximum for orbit 170. The model magnetic field profile tends to be sensitive to the solar wind temperature, and because that is known to within only a factor of two, the differences between the computed and measured magnetic field profiles are reasonable. The
particle pressure drops in this region as well, but undergoes a resurgence as the flow moves into the oxygen ledge enhancement region. The computed magnetic field pressure drops throughout the ledge region to balance the total pressure. This is contrary to the observed magnetic field profile which appears to maintain a high value through the oxygen ledge region. The calculated temperatures (Figure 15) in the ledge region are a factor of 10 higher than those measured by the OETP (orbital electron temperature probe) instrument for orbit 170. This may explain the rapid resurgence of the particle pressure in this region. Were the temperature to be lower as in the observations, the particle pressure would remain low and the magnetic field would continue to dominate throughout this region as it does observationally.

Computer simulations are compared to the measured ionopause regions for various orbits in Figures 16-18. In general, the calculations exhibit similar scale heights to the observed profiles, although the differences in the measured altitudes may be as much as 100 km. Orbit 188 shows particularly good correspondence with the measured ionopause profile. For orbit 184, the OMAG instrument gives a magnetic field maximum of 91 nT, while the calculations give a 77 nT maximum. Similarly, for orbit 408, the measured maximum is 104 nT and the calculated is 114 nT, and for orbit 188, 88 nT is the measured maximum where 85 nT is calculated. It is apparent that the correspondence between the calculated and observed magnetic field maxima is close.

The shock distances in the calculations needed to create reasonable velocity profiles tend to be significantly higher than those observed by the
OMAG instrument\textsuperscript{40}. Since, to a certain extent, the shock distances are an artificial parameter to give the proper velocity ratios, their accuracy must be considered suspect independent of observations. To compute a proper shock shape and radius, a full three dimensional flow model would have to be constructed. For the present model, it will suffice to say that mass-loading appears to increase the shock radius, at least at Venus, in agreement with the results of Breus et al.\textsuperscript{41}

5.3 CONCLUSIONS

Whether the oxygen density ledge computed by this model may be the ionopause or not is still difficult to determine at this time. The ledge shows the same scale height and density features as the observed structure, yet it does not simulate the dynamic response of the ionopause, especially for low solar wind pressures where the observed ionopause may be hundreds of kilometers higher than the computed ledge. On a positive note, it should be stated that this model is the best one produced to date. No other model has attempted to reconstruct the effects of the magnetic field on the mantle flow structure, despite universal consent as to the importance of the magnetic pressure above the ionopause. It would seem that a full understanding of the mass-loaded "Zwan-Wolf-like" effect should be critical to a complete description of the solar wind-Venus interaction. One of the purposes of this research is to make positive contributions to the formulation of such a complete picture.

Obvious areas of future work include the expansion of the Ohm's Law to include those terms which may be most important in regions where
collisions dominate. The expansion of the plasma velocity parameters to allow asymmetric flow would also be useful in creating a more general formulation of the mass-loaded Zwan-Wolf effect. The resulting density effects may solve some of the problems now associated with the dynamic behavior of the ionopause. Another obvious future goal would be to include the effects of non-thermal species. The extension of the model to include higher order terms than those of equations (44a) - (44e) might also improve the model calculations.

Despite the relatively long history of Venus studies, there is still much disagreement over explanations of many of the basic phenomena. Considering the potentially widespread nature of plasma-neutral gas interactions throughout the universe, a proper understanding of the Venus interaction is paramount. I hope that this work has made a positive contribution in this area.


40. T. Zhang, UCLA, private communications.

Figure 1

Diagram of the solar wind interaction regions of Venus. The ionopause marks the separation between the mantle and the ionosphere.
Diagram of ion pickup in a flowing magnetized plasma. The plasma is flowing in the $x$ direction and the magnetic field points in the $z$ direction. (a) From the rest frame of the ion, the plasma creates a $-v \times B$ electric field which, in conjunction with the magnetic field, drags the ion along with the rest of the plasma. (b) From the rest frame of the plasma, the ion is added at velocity $-v$ and subsequently gyrates in standard cyclic motion with the guiding center moving at the same velocity as the plasma.
Figure 3

Diagram of planetocentric spherical coordinates used in the model. $r$ is the radius measured from the center of the planet. The stagnation flow line lies along $\theta = 0$. 
The magnetic field strengths from the Lees model are here compared to the present model. In this graph and those that follow, the plasma flow parameters are normalized to their post-shock values. The altitude is also normalized so that 1.0 is the shock altitude and 0.0 is the body surface. The magnetic profiles match closely near the shock, but the present model shows a stronger field modification near the body surface.
The ion density for the Lees model monotonically approaches a zero value as the plasma moves toward the body surface, while the present model shows a slight increase behind the shock before approaching zero near the body surface.
The downward velocities of both the current model and the Lees model decrease monotonically to zero near the body surface.
Comparison of Magnetic Fields for Spreiter Model

Figure 7

Comparison of the current model to the Spreiter model shows good agreement for the magnetic field profiles. Both models show the magnetic field strength increasing to large values near the body surface.
Comparison of Densities for Spreiter Model

Figure 8

The computed ion density of the Spreiter model shows a similar profile to that of the present model, but not as large a modification near the body surface. Note that the differences between the two models are only a few per cent of the total density.
As with the Lees model, the Spreiter model and the present model show similar downward velocity profiles.
The plasma temperature profiles for both the Spreiter model and the present model are similar to those of density (Figure 8). As before, the difference between the present model and that of Spreiter is only a few per cent.
Using measured solar wind conditions, I have calculated the predicted velocity profiles for orbit 170. The v profile (recall that $v \cdot r = v \cos \theta$) is shown with downward velocity as positive. The downward velocity drops asymptotically from its maximum of 1.5 km/sec behind the shock at the altitude of 3200 km to a value predicted to be only a few meters per second near the ionosphere. Note that the u profile (recall that $v \cdot \theta = v \sin \theta$) maintains a nearly constant value except near the ionosphere where it drops toward zero; possibly indicating a transition toward more vertical flow of the plasma through the ionopause.
Computed Ion Concentrations for Orbit 170

As can easily be seen from this graph for orbit 170, hydrogen dominates the ion density through most of the mantle except near the ionosphere, where the oxygen density dominates. Just above the ionopause region, near 400 km, there is a region of plasma depletion corresponding to an increase in the magnetic field pressure (see Figure 14). This region also corresponds to a slight increase in u (see Figure 11), indicating that much of the plasma is diverted around the planet in this region.
This is a magnification of the lower section of Figure 12 showing the onset of the oxygen "ledge". Also plotted are ion densities in the ionopause region for orbit 170 measured by the Orbital Ion Mass Spectrometer (OIMS) aboard the Pioneer Venus spacecraft. While the scale heights of the oxygen "ledge" and the ionopause correspond well, the two structures appear separated by approximately 75 km. The differences between the measured ionopause and the oxygen ledge are due to the inability of the present model to recreate the dynamic behavior that the ionopause shows in response to different solar wind parameters.
The computed magnetic field strength for orbit 170 increases to a maximum in the same region where the ion density dips (see Figure 12). This may indicate a transition from plasma pressure dominated flow to one of magnetic pressure dominated flow. Note the drop in the magnetic field through the oxygen "ledge" region. Such a drop in the field strength is not seen at Venus for this orbit. The model may overestimate the plasma pressure in this region, underestimating the magnetic pressure to balance the total pressure.
The computed plasma temperatures for orbit 170 show a very hot plasma through much of the mantle. The temperatures near the ionopause appear to be higher than those measured at Venus by a factor of 10, perhaps explaining the magnetic field strength drop near the ionopause (see Figure 14).
The computed ion densities for orbit 184 resemble those of orbit 170. The
distance between the actual ionopause measured by the OIMS instrument
and the computed oxygen "ledge" is approximately 100 km. As with orbit
170, the scale heights of the two structures are similar.
Figure 17

The ionopause measured by the OIMS instrument for orbit 188 corresponds quite well with the computed oxygen "ledge".
Figure 18

Orbit 408 shows a separation between the ionopause and the computed oxygen "ledge" of approximately 100 km, much like orbit 184 (Figure 16) and orbit 170 (Figure 13).