# Ionosphere/thermosphere heating determined from dynamic magnetosphere-ionosphere/thermosphere coupling

Jiannan Tu,<sup>1</sup> Paul Song,<sup>1</sup> and Vytenis M. Vasyliūnas<sup>1,2</sup>

Received 4 March 2011; revised 7 August 2011; accepted 16 August 2011; published 17 September 2011.

[1] Ionosphere/thermosphere heating driven by magnetospheric convection is investigated through a three-fluid inductive (including Faraday's law) approach to describing magnetosphere-ionosphere/thermosphere coupling, for a 1-D stratified ionosphere/ thermosphere in this initial study. It is shown that the response of the ionosphere/ thermosphere and thus the heating is dynamic and height-dependent. The heating is essentially frictional in nature rather than Joule heating as commonly assumed. The heating rate reaches a quasi-steady state after about 25 Alfvén travel times. During the dynamic period, the heating can be enhanced and displays peaks at multiple times due to wave reflections. The dynamic heating rate can be more than twice greater than the quasi-steady state value. The heating is strongest in the E-layer but the heating rate per unit mass is concentrated around the F-layer peak height. This implies a potential mechanism of driving O<sup>+</sup> upflow from O<sup>+</sup> rich F-layer. It is shown that the ionosphere/ thermosphere heating caused by the magnetosphere-ionosphere coupling can be simply evaluated through the relative velocity between the plasma and neutrals without invoking field-aligned currents, ionospheric conductance, and electric field. The present study provides understanding of the dynamic magnetosphere-ionosphere/thermosphere coupling from the ionospheric/thermospheric view in addition to magnetospheric perspectives.

Citation: Tu, J., P. Song, and V. M. Vasyliūnas (2011), Ionosphere/thermosphere heating determined from dynamic magnetosphere-ionosphere/thermosphere coupling, *J. Geophys. Res.*, *116*, A09311, doi:10.1029/2011JA016620.

## 1. Introduction

[2] The ionosphere/thermosphere heating is an essential process that has a strong influence on many important phenomena. For example, the increase of thermospheric scale height during magnetospheric storms can remarkably increase the satellite drag and thus affect satellites on the low-altitude orbits. Therefore, it is crucial to understand the ionosphere/thermosphere heating and determine the heating rate during the dynamic periods of the magnetosphere-ionosphere/thermosphere system.

[3] The heating caused by the electromagnetic energy input associated with the solar wind-magnetosphere-ionosphere/ thermosphere coupling is conventionally represented by the Joule heating  $\mathbf{J} \cdot (\mathbf{E} + \mathbf{U} \times \mathbf{B})$  [e.g., *Kelley*, 1989; *Richmond*, 1995], where **B**, **U**, **E** and **J** is the geomagnetic field, neutral wind velocity, electric field in the Earth frame, and electric current density, respectively. Since **E** and **U** depend on the frame of reference, they must be measured in the same frame. The implementation of the Joule heating calculation is commonly based on quasi-steady state Ohm's law, in which the conductivities are real [e.g., *Richmond et al.*,

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1992; *Fuller-Rowell et al.*, 1996; *Ridley et al.*, 2006]. Under the quasi-steady state assumption, the inertial terms in electron and ion momentum equations are neglected and the electric field is assumed curl-free. The response of the ionosphere/ thermosphere to magnetospheric perturbations, however, does not instantly reach the quasi-steady state. There exists a transient period when the response is dynamic (non-negligible time derivatives) and the quasi-steady state Joule heating description in the conventional theories may not be applicable.

[4] The question here is, how long is the transient period? Let us consider perturbations that are induced by the solar wind-magnetosphere interactions at the magnetopause and propagate to the ionosphere along the magnetic field lines with the Alfvén velocity. The perturbations are partially reflected from the ionosphere due to the density gradient there. The entire magnetosphere-ionosphere/thermosphere (M-IT) system may reach a steady state after a number of back and forth bounces of the perturbations [e.g., Lysak and Dum, 1983]. There are three timescales involved in this process. The shortest one is the Alfvén travel (or transient) time  $t_A$ , which is the time for the perturbations to propagate from the magnetopause to the ionosphere. Assuming the magnetopause to be 15  $R_E$  away from the polar ionosphere and with an average Alfvén velocity of 1000 km/s, we obtain  $t_A \sim 100$  s. The longest timescale is that for a significant change of the bulk velocity of the neutrals, which is also the timescale to reach a steady state (if one exists) of the entire M-IT system, establishing equilibrium among all three subsystems (the magnetosphere, the ionosphere, and the neutral thermosphere).

<sup>&</sup>lt;sup>1</sup>Center for Atmospheric Research, University of Massachusetts Lowell, Lowell, Massachusetts, USA.

<sup>&</sup>lt;sup>2</sup>Max-Planck-Institut für Sonnensystemforschung, Katlenburg-Lindau, Germany.

This timescale is of the order of  $1/\nu_{ni}$ , where  $\nu_{ni}$  is the neutral-ion collision frequency, and is about 1–3 hours in the F layer of the Earth's ionosphere but can be much longer at lower altitudes. The third, intermediate timescale, for establishing flow equilibrium between the magnetosphere and the ionosphere (but not the thermosphere), is typically of the order of 15–20 min, as shown in some previous studies [e.g., *Holzer and Reid*, 1975; *Vasyliūnas and Pontius*, 2007; *Song et al.*, 2009]. This timescale is in the range of or longer than the timescales of many important dynamical phenomena in the ionosphere/thermosphere. Therefore, the heating and in general the solar wind-magnetosphere-ionosphere/thermosphere coupling in the transient period need to be properly considered.

[5] The dynamic magnetosphere-ionosphere coupling has been investigated in numerous studies [e.g. Lighthill, 1960; Hughes, 1974; Holzer and Reid, 1975; Hughes and Southwood, 1976; Newton et al., 1978; Lysak and Dum, 1983; Allen et al., 1987; Kivelson and Southwood, 1988; Wright, 1996; Lysak, 1999, 2004; Sciffer and Waters, 2002; Wolf et al., 2006]. For instance, Lysak and Dum [1983] investigated dynamic coupling of the magnetosphere with the ionosphere under various types of specified sources at the equatorial plane. They found that the magnetosphereionosphere system would experience a number of back and fourth bounces of perturbations between the ionosphere and source region to reach the steady state. In those studies, the ionosphere is treated as a height-integrated lower boundary of the magnetosphere at which waves reflect [e.g., Lysak and Dum, 1983; Wright, 1996]. In the height-integrated ionosphere neither the wave reflection nor the plasma and neutral dynamics within the ionosphere are resolved. Hughes [1974] and Hughes and Southwood [1976] studied the screening effect of the atmosphere and ionosphere on low-frequency hydromagnetic waves. Their approach includes the inductive effects as well as plasma dynamics in a structured ionosphere and is in principle able to describe the dynamics of the M-IT coupling, but it has not yet been developed into a full-fledged explicit description of the plasma and neutral dynamics of the M-IT system.

[6] In the present study we investigate the heating of the ionosphere/thermosphere driven by the magnetospheric convection on the basis of a self-consistent treatment of plasma and neutral dynamics and electromagnetic fields in describing the ionosphere/thermosphere and its coupling with the magnetosphere. The basic equations and the numerical scheme are described in section 2; simulation results are presented in section 3; discussion and conclusions are given in section 4.

### 2. Basic Equations and Numerical Scheme

[7] The ionosphere/thermosphere can be thought to consist of three fluids, after lumping all ion species into one ion fluid and all different neutrals into one neutral fluid: electron, ion, and neutral fluids [e.g., *Richmond*, 1995; *Song et al.*, 2005a]. The governing equations for this system are the generalized Ohm's law, plasma and neutral momentum equations, plus Maxwell equations [*Song et al.*, 2005a, 2005b]

$$\frac{-m_e}{e}\frac{\partial \mathbf{J}}{\partial t} = \mathbf{J} \times \mathbf{B} - eN_e(\mathbf{E} + \mathbf{V} \times \mathbf{B}) + \frac{m_e}{e}(\nu_{en} + \nu_{ei})\mathbf{J} - N_e m_e(\nu_{en} - \nu_{in})(\mathbf{V} - \mathbf{U})$$
(1)

$$N_e m_i \frac{\partial \mathbf{V}}{\partial t} = \mathbf{J} \times \mathbf{B} - N_e (m_i \nu_{in} + m_e \nu_{en}) (\mathbf{V} - \mathbf{U}) + \frac{m_e}{e} (\nu_{en} - \nu_{in}) \mathbf{J}$$
(2)

$$N_n m_n \frac{\partial \mathbf{U}}{\partial t} = N_e (m_i \nu_{in} + m_e \nu_{en}) (\mathbf{V} - \mathbf{U}) - \frac{m_e}{e} (\nu_{en} - \nu_{in}) \mathbf{J} \quad (3)$$

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E}, \quad \mu_0 \mathbf{J} = \nabla \times \mathbf{B}$$
(4)

where V is the plasma bulk velocity;  $N_e$  and  $N_n$  is the electron and neutral number density, respectively;  $m_e$ ,  $m_i$ , and  $m_n$  is the electron mass, mean ion mass, and mean neutral mass, respectively; e is the elementary charge;  $\nu_{en}$  and  $\nu_{ei}$  is the electron-neutral and electron-ion collision frequency, respectively; and  $\mu_0$  is the permeability in vacuum. The equations have been written assuming charge quasi-neutrality. The continuity equations that determine the densities are not shown because for the simple geometry assumed in this initial study they are automatically satisfied, as noted below, and the densities can be treated as constant in time. Note that above equations do not include the conventional ionospheric Ohm's law, which holds only under the assumption of a steady state. Plasma electrodynamics are treated self-consistently. Neutral dynamics are included self-consistently as far as the effect of collisions with the plasma are concerned; other forces on neutrals are, as a first approximation, neglected or else assumed to be balanced among themselves.

[8] In this initial study we consider a 1-D stratified ionosphere/thermosphere: spatial variation is only along the vertical magnetic field (constant background magnetic field  $\mathbf{B} = -B_0 \hat{z}$  with  $\hat{z}$  vertically upward). Furthermore we neglect vertical flow and current (those effects will be included in a future model development). For this simplified system, the divergences of the flows are zero and hence the continuity equations are automatically satisfied.

[9] The energy input to the ionosphere-thermosphere contains dissipation or thermal heating. The total heating rate, including both the ionosphere (plasma) and the thermosphere (neutrals), is given by *Vasyliūnas and Song* [2005] as

$$q = \mathbf{J} \cdot (\mathbf{E} + \mathbf{V} \times \mathbf{B}) + \rho \nu_{in} (\mathbf{V} - \mathbf{U})^2 = \eta \mathbf{J}^2 + \rho \nu_{in} (\mathbf{V} - \mathbf{U})^2 \quad (5)$$

where  $\rho = N_e m_i$  is the plasma mass density and  $\eta = m_e (\nu_{en} + \mu_{en})$  $\nu_{ei}/e^2 N_e$  is the Ohmic resistivity. We have applied the generalized Ohm's law (the time derivative of J, which is much smaller compared to the other terms [Vasyliūnas and Song, 2005], particularly when charge quasi-neutrality holds [Vasyliūnas, 2005a, 2005b, 2011], has been neglected) in order to obtain the second expression of (5). The term  ${\bf J}$  ·  $(\mathbf{E} + \mathbf{V} \times \mathbf{B})$ , associated with the electric field in the plasma frame of reference, is true Joule heating because it is the same as that in the original form of Ohmic heating  $\eta \mathbf{J}^2$ . The term  $\rho \nu_{in} (\mathbf{V} - \mathbf{U})^2$  is the frictional heating caused by the collisions between the plasma and neutrals. It should be pointed out that in this study we calculate the heating rates but not the explicit changes of pressure or the effects on the solutions of equations (1)-(4), which will be considered in a future study by incorporating the energy equations.



**Figure 1.** Altitude profiles of neutral and plasma density, ion-neutral and neutral-ion collision frequencies.

[10] Using Ampère's law to substitute current density in (1)–(3) and using (1) to replace electric field in Faraday's law, one obtains a set of equations for the plasma and neutral velocity V and U, and perturbation magnetic field  $\mathbf{B}_{\perp}$ , which for the assumed 1-D geometry is perpendicular to the ambient magnetic field. This equation set describes the dynamics of the plasma, neutrals, and magnetic field on MHD timescales for the simple 1-D stratified ionosphere/thermosphere and is numerically solved using a fully implicit difference method. The resulting difference equations with normalized variables are

$$\overline{\mathbf{V}}_{j}^{n+1} + \overline{\nu}_{in,j}\Delta\overline{t}\left(\overline{\mathbf{V}}_{j}^{n+1} - \overline{\mathbf{U}}_{j}^{n+1}\right) + \frac{\Delta\overline{t}\left(\overline{\mathbf{B}}_{\perp,j+1}^{n+1} - \overline{\mathbf{B}}_{\perp,j-1}^{n+1}\right)}{2\overline{\rho}_{j}\Delta\overline{z}} - \frac{\left(\overline{\nu}_{en,j} - \overline{\nu}_{in,j}\right)\Delta\overline{t}}{2\overline{\rho}_{j}\overline{\Omega}_{e0}\Delta\overline{z}}\hat{\mathbf{z}} \times \left(\overline{\mathbf{B}}_{\perp,j+1}^{n+1} - \overline{\mathbf{B}}_{\perp,j-1}^{n-1}\right) = \overline{\mathbf{V}}_{j}^{n}$$
(6)

$$\overline{\mathbf{U}}_{j}^{n+1} - \alpha_{j}\overline{\nu}_{in,j}\Delta\overline{t}\left(\overline{\mathbf{V}}_{j}^{n+1} - \overline{\mathbf{U}}_{j}^{n+1}\right) + \frac{\alpha_{j}\left(\overline{\nu}_{en,j} - \overline{\nu}_{in,j}\right)\Delta\overline{t}}{2\overline{\rho}_{j}\overline{\Omega}_{e0}\Delta\overline{z}}\hat{z}$$
$$\times \left(\overline{\mathbf{B}}_{\perp,j+1}^{n+1} - \overline{\mathbf{B}}_{\perp,j-1}^{n+1}\right) = \overline{\mathbf{U}}_{j}^{n} \tag{7}$$

$$\overline{\mathbf{B}}_{\perp,j}^{n+1} + \frac{\Delta \overline{t}}{2\Delta \overline{z}} \left( \overline{\mathbf{V}}_{j+1}^{n+1} - \overline{\mathbf{V}}_{j-1}^{n+1} \right) - \frac{\Delta \overline{t}}{4 \overline{\Omega}_{i0} \Delta \overline{z}^2} \left[ \left( \beta_{j+1} - \beta_{j-1} \right) \right. \\ \left. \cdot \left( \overline{\mathbf{B}}_{\perp,j+1}^{n+1} - \overline{\mathbf{B}}_{\perp,j-1}^{n+1} \right) + 4 \beta_j \left( \overline{\mathbf{B}}_{\perp,j+1}^{n+1} - 2 \overline{\mathbf{B}}_{\perp,j}^{n+1} + \overline{\mathbf{B}}_{\perp,j-1}^{n+1} \right) \right] \\ \left. - \frac{\Delta \overline{t}}{\overline{\Omega}_{i0} \overline{N}_{e,j} \Delta \overline{z}^2} \hat{\mathbf{z}} \times \left[ \left( \overline{\mathbf{B}}_{\perp,j+1}^{n+1} - 2 \overline{\mathbf{B}}_{\perp,j}^{n+1} + \overline{\mathbf{B}}_{\perp,j-1}^{n+1} \right) \right] \\ \left. - \frac{\left( \overline{N}_{e,j+1} - \overline{N}_{e,j-1} \right)}{4 \overline{N}_{e,j}} \left( \overline{\mathbf{B}}_{\perp,j+1}^{n+1} - \overline{\mathbf{B}}_{\perp,j-1}^{n+1} \right) \right] = \overline{\mathbf{B}}_{\perp,j}^{n} \tag{8}$$

where the superscript n + 1 represents the (n + 1)th time step and the subscript *j* is the spatial grid index,  $\Delta \overline{t}$  and  $\Delta \overline{z}$  is the normalized time step and spatial grid interval, respectively,  $\overline{\rho}$  is the normalized plasma mass density,  $\alpha = N_e m_i / N_n m_n$ ,  $\overline{\Omega}_{e0}$  and  $\overline{\Omega}_{i0}$  is the normalized electron and ion gyrofrequency, respectively, at the top boundary, and  $\beta = (\overline{\nu}_{en} + \overline{\nu}_{ei})/\overline{\Omega}_{e0}\overline{N}_e$ . The normalization is that  $\overline{\mathbf{V}} = \mathbf{V}/V_{A0}$ ,  $\overline{\mathbf{U}} = \mathbf{U}/V_{A0}$ ,  $\overline{\mathbf{B}}_{\perp} = \mathbf{B}_{\perp}/B_0$ ,  $\overline{N}_e = N_e/N_{e0}$ ,  $\overline{t} = t/t_0$ ,  $\overline{z} = z/L$ ,  $\overline{\Omega}_{e0} = \Omega_{e0}t_0$ ,  $\overline{\Omega}_{i0} = \Omega_{i0}t_0$ ,  $\overline{\nu}_{en} = \nu_{en}t_0$ ,  $\overline{\nu}_{in} = \nu_{in}t_0$ , and  $\overline{\nu}_{ei} = \nu_{ei}t_0$ , where  $N_{e0}$  and  $V_{A0}$  is the electron density and Alfven velocity, respectively, at the top boundary, and  $t_0 = L/V_{A0}$  with L the length of the simulation domain.

[11] The current density J is calculated from Ampère's law and the electric field E from the generalized Ohm's law, using the solution from (6)-(8). The modeled ionosphere/ thermosphere system is assumed to be driven by a convection velocity at the top boundary (1000 km). The antisunward component  $V_x$  changes from 0 to  $0.001 V_{A0}$  ( $V_{A0} \approx$ 3000 km/s in the present case) in 1 s. The initial magnetic field is  $\mathbf{B} = -B_0 \hat{z}$  with  $B_0 = 50000$  nT. The initial (or background) plasma and neutral velocities are set to zero. At the top boundary, the spatial derivative of the perturbation magnetic field (or, equivalently, the current density) is assumed to be zero. At the bottom boundary (80 km), for the magnetic field we likewise set the perturbation magnetic field to zero, and we assume that the spatial derivative of the plasma velocity is zero. The spatial resolution of the calculation is 4 km and the time step is 0.1 s. Other time steps, such as 1 s and 0.01 s, have been tested, and no distinguishable difference between the results with different time steps is found.

[12] The implicit difference equations (6)–(8) are unconditionally stable and can be solved using a much longer time step than with explicit methods. To solve (1)–(4) [Song et al., 2009] used an explicit difference method, which requires a very small time step  $(10^{-7} \text{ s})$ . A 6-powers-of-10 longer time step of the implicit scheme adopted here greatly increases the calculation speed, even though the implicit differencing leads to a set of linear algebraic equations with a large matrix. The results from implicit and explicit methods are essentially the same, as will be shown later in Figure 3, providing us with confidence that the numerical solutions are reliable.

[13] The background ionospheric parameters, including the electron density, the mean ion mass  $m_i$ , and electron and ion temperatures (used to calculate collision frequencies), are taken from the International Reference Ionosphere (IRI) 2007 [Bilitza and Reinisch, 2008]. The calculation is for the northern pole in winter solstice and at local magnetic noon. The background neutral density, the mean neutral mass, and the neutral temperature are evaluated using the NRLMSISE00 empirical atmospheric model [Picone et al., 2002]. The ionneutral, electron-neutral, and electron-ion collision frequencies are evaluated using the formulas given by Schunk and Nagy [2000]. Figure 1 displays the altitude profiles of the neutral and the plasma densities as well as the ion-neutral and the neutral-ion collision frequencies  $\nu_{in}$  and  $\nu_{ni}$ . The neutral density is much higher than the plasma density below 600 km, decreases exponentially with increasing altitude, and becomes smaller than the plasma density above 800 km. The ion-neutral collision frequency is 10<sup>6</sup> Hz at 80 km but quickly decreases to less than 0.01 Hz at altitudes above 600 km. The neutral-ion collision frequency shows an altitude dependence similar to that of the plasma (ion) density, with a maximum in the F-layer. This altitude distribution of  $\nu_{ni}$ , the result of the momentum conservation during collisions  $\rho_n \nu_{ni} = \rho \nu_{in}$ , is the cause of distinct long time varia-



**Figure 2.** The relative velocity normalized to  $V_{A0}$ , with (a) x direction antisunward and (b) y dawnward, as a function of altitude for the first 40 s after the top boundary starts moving.

tions of the heating rate at different altitudes, as explained in section 3.2.

## 3. Convection-Driven Frictional Heating

## 3.1. Plasma and Neutral Motion

[14] Before discussing the heating rate we first describe briefly the behavior of the plasma and neutrals in response to an imposed convection velocity at the top boundary, for the purpose of understanding the time variation and altitude distribution of the heating rate. Let us consider the process in a simple 1-D stratified ionosphere/thermosphere with a vertical magnetic field. The plasma and neutrals are initially (t = 0) at rest (or at the same background velocity). A plasma motion is imposed at the top boundary at t = 0+. The plasma inside the domain is linked by the magnetic field. The plasma motion creates a kink on the magnetic field line at the top boundary because of the frozen-in condition of the magnetic field and the plasma. The kink exerts a tension force on the plasma below the boundary and makes it move, creating another kink at lower altitude, and so on. The net result is a magnetic perturbation front propagating downward along the magnetic field line at the Alfvén speed  $V_A$ . This is the primary mechanism of momentum coupling from the magnetosphere to the ionosphere. The neutrals, on the other hand, are not affected by the electromagnetic force. After the perturbation front has passed, there is a velocity difference between the plasma and neutrals. The relative motion causes, through ion-neutral collisions, momentum transfer between plasma and neutrals and frictional heating of both. The collisions tend to slow down the plasma, but the net effect is barely noticeable on timescales shorter than  $1/\nu_{in}$ , and the plasma motion continues at all times as long as the motion at the top boundary is maintained. Meanwhile, the neutrals are continually accelerated by the collisions and after times  $t \sim 1/\nu_{in}$  are getting close to catching up with the plasma. In the absence of other forces on the neutrals, the system will eventually, when t goes to infinity, reach a steady state in which the plasma and the neutrals have a common velocity equal to the driving velocity.

[15] The above discussion applies only to one-way propagation of the perturbations, but it does provide a basic understanding of how the plasma and the neutrals respond to the driver in a time-dependent fashion. In reality the density gradients of the ionosphere/thermosphere (and to a lesser degree the magnetic field gradient, which is not considered here because of the assumed uniform background magnetic field in our model) cause partial reflection of the downward propagating perturbation, modifying the basic picture discussed above. Wave propagation and reflection are therefore key processes in the magnetosphere-ionosphere/ thermosphere coupling [e.g., *Lighthill*, 1960; *Hughes*, 1974; *Lysak and Dum*, 1983; *Sciffer and Waters*, 2002]. Numerical solutions of the equation set (1)–(4) automatically include the reflections (without the need to specify explicitly a reflection coefficient) and thus provide a quantitative description of the modified process. Since our focus is on the heating which, as will be shown later, is essentially frictional, it suffices to examine only the relative velocity V - U obtained from the numerical solutions.

[16] Figure 2 displays the altitude distribution of the relative velocity for the first 40 s after the top boundary starts moving. During this early period, the neutrals have not gained any noticeable motion so that the relative velocity is essentially the plasma velocity. The topside velocity transition occurs in 1 s as assumed. The initial perturbation reaches the bottom boundary in about 1 s or close to an Alfvén travel time. Hereafter, when discussing wave reflection and transient processes, the time is given in units of the Alfvén travel time, defined as the integrated time for the perturbation to propagate from the top to the bottom boundary, taking into account the density variation with altitude and using the local Alfvén velocity. The actual propagation time is slightly (about 10%) longer than the nominal Alfvén travel time calculated from the local Alfvén speed because of a neutralinertia loading effect associated with heavy collisions at lowest-altitudes of the ionosphere [Song et al., 2005b]. Representing the time in units of the Alfvén travel time is useful for scaling. In the present simulation, with upper boundary at 1000 km. the Alfvén travel time is :0.9 s but it will be of order  $10^2$  s if the upper boundary of the simulation domain is placed at a realistic distance (say, 15  $R_e$  from the polar cap to the magnetopause).

[17] It is seen from Figure 2 that the full strength of the imposed convection at the top boundary is felt at the bottom after about two Alfvén travel times. The relative velocity (approximately the plasma velocity) shows large variations in response to the driver during the first 10 Alfvén travel times, and the antisunward velocity profile reaches an overall quasi-steady state in about 25 Alfvén travel times. While the



**Figure 3.** Altitude variation of the frictional heating rate (solid line) and the true Joule heating rate (dashed line) at 10 Alfvén travel times. The dotted line is the heating rate from the simulation with FTCS difference scheme.

plasma motion has settled down, the neutrals continue to be accelerated through the ion-neutral collisions. After a time  $1 \sim /\nu_{ni}$  the neutrals have acquired a significant velocity and the relative velocity slowly decreases (simulation results not shown).

[18] Figure 2 also shows that the imposed antisunward motion ( $V_x$  component) at the top boundary not only drives antisunward motion at lower altitudes but also induces a dawnward (y) component, caused by the Hall effect. The x component of the magnetic tension force also causes an initial overshoot in the antisunward velocity for a short time period (thus producing the strongest heating, as will be shown later). After this temporary overshoot, the relative velocity in the E-layer is small, but large velocity values (up to the driving velocity) are seen above 140 km, i.e., in the F-layer and topside ionosphere.

#### 3.2. Heating Rate

[19] Now we examine the heating rate calculated from the numerical solutions of equations (1)–(4). We first compare the two terms on the right hand side of equation (5): Joule or Ohmic heating rate  $\mathbf{J} \cdot (\mathbf{E} + \mathbf{V} \times \mathbf{B}) = \eta \mathbf{J}^2$  and frictional

heating rate  $\rho \nu_{in} (\mathbf{V} - \mathbf{U})^2$ . Displayed in Figure 3 are altitude variations of Joule and frictional heating rates at a given time. It is clear from Figure 3 that the frictional heating is dominant except at lowest ionospheric altitudes. Only at altitudes close to 80 km is the Joule heating larger than the frictional, but with values about two orders of magnitude smaller than the peak frictional heating rate. The height-integrated Joule heating rate is negligibly small at all times, indicating that heating is primarily frictional.

[20] In Figure 3 we also show the heating rate (dotted line) calculated from the simulation using an explicit difference scheme, namely, the forward-time-centered-space (FTCS) method [e.g., *Morton and Mayers*, 2005]. The difference in the frictional heating rates calculated from the implicit and explicit methods is so small that they overlap as one solid line. The difference in the Ohmic heating rates (dashed and dotted lines) is noticeable but very small. In addition, no significant differences in plasma and neutral velocities and magnetic field are found between two schemes (results not shown). The fact that the results from two quite different numerical schemes are essentially the same is evidence that our numerical solutions are reliable.

[21] Having shown that the heating is essentially frictional, we can understand the time variation and spatial distribution of the heating rate from the behavior of the relative motion. In Figure 4 we present a contour plot of heating rate q versus time (in units of Alfvén travel time) and altitude. Large variations in the heating rate are seen during the first 10 Alfvén travel times, in association with the large variations in the relative velocity (note that the heating rate is positive for both positive and negative velocity differences). The strongest heating occurs when  $V_x$  overshoots the driving convection velocity, due to wave reflection. As the relative motion approaches a quasi-steady state after about 25 Alfvén travel times, so does the heating rate.

[22] The altitude distribution of the heating rate is controlled by three factors: plasma mass density  $\rho$ , ion-neutral collision frequency  $\nu_{in}$ , and the square of the relative velocity (V – U). The relative velocity is small below about 100 km. The plasma mass density is maximum at the F-layer peak height but is small in the E-layer. The collision frequency  $\nu_{in}$ , however, increases monotonically with decreasing altitude:



Figure 4. Heating rate q as a function of Alfvén travel time and altitude.



**Figure 5.** Heating rate per unit mass (in unit of erg  $g^{-1} s^{-1}$ ),  $q/\rho_t$ , where  $\rho_t$  is the total mass density, as function of Alfvén travel time and altitude.

it is of the order of  $10^6$  Hz at 80 km and about 100 Hz at 120 km (see Figure 1). The strongest heating concentrated in the E-layer around 120 km can be understood as the result of the relatively large ion-neutral collision frequency there. Farther down toward the bottom of the ionosphere, the collision frequency is even larger but the relative velocity and the plasma density are extremely small. A secondary peak of the heating rate occurs in the F layer due to the plasma mass density maximum around the F-layer peak height and larger relative velocity. At higher altitudes, the heating rate decreases rapidly because of the exponentially decreasing ion-neutral collision frequency, even though the relative velocity remains high.

[23] Figures 3 and 4 show that the heating is strongest in the E-layer, but the total number density of the particles there is also high. The effectivenesses of the heating can be evaluated by dividing the heating rate q by the total mass density (plasma plus neutral mass density) to obtain the heating rate per unit mass (equivalent to the heating per particle), shown in Figure 5. The heating per unit mass is actually largest in the F-layer, around 350 km, and becomes weak below about 250 km because of the large neutral mass density there. This result can be understood by noting that the relation  $\rho_n \nu_{ni} = \rho_i \nu_{in}$  (momentum conservation during collisions) inserted into the frictional heating rate (5) gives heating rate per unit mass  $\nu_{ni}(\mathbf{V} - \mathbf{U})^2$ , directly proportional to the neutral-ion collision frequency. The frequency  $\nu_{in}$ , as shown in Figure 1, is peaked at the F-layer and becomes small below about 250 km. The concentrated heating rate per unit mass means that the energy received by each particle is largest in the F-layer. Because the F-layer is rich in O<sup>+</sup>, the large heating rate per unit mass has significant consequences for heavy ion (O<sup>+</sup>) outflow and may provide an explanation for the enhanced  $O^+$  ion outflow in response to a solar wind pressure pulse [Fuselier et al., 2002] or an interplanetary magnetic cloud [Zong et al., 2008].

[24] Since the frictional heating rate is proportional to the plasma mass density, equal to  $N_e m_i$ , the electron density  $N_e$  directly influences the heating rate. In the present study the  $N_e$  density profile deep inside the polar cap is used, with  $N_e$ 

 $\sim 4 \times 10^5$  cm<sup>-3</sup> at the F<sub>2</sub> peak height ( $N_m F_2$ ). With a larger  $N_e$ , e.g., in the cusp, the calculated heating rate may be significantly greater. More importantly, the plasma velocity may be larger around the cusp, leading to localized high heating rates there, which may be one of causes for the enhanced neutral mass density observed by the CHAMP satellite over the cusp [*Luhr et al.*, 2004].

[25] As shown in Figure 4, the strongest heating occurs during the transient period. This is more clearly apparent in the height-integrated heating rate displayed in Figure 6. The segment highlighted with diamonds is the heating rate after the quasi-steady state has been reached. It is seen that during the first 20 Alfvén travel times the heating rate is in general larger than in the quasi-steady state, with peak heating rate double the quasi-steady state value. Note that heating is not significant until the perturbation has arrived in the high collision frequency region. The peaks in the heating rate during the dynamic period demonstrate that wave reflection plays a key role in producing enhanced heating: the reflected perturbation may either enhance or reduce the strength of the incident perturbation, depending on the



Figure 6. Time variation of height-integrated heating rate.



Figure 7. Same format as in Figure 4 but for an extended time period.

phase delay. The heating rate first displays two peaks separated by about two Alfvén travel times, indicating that during the first two bounce periods the reflection enhances the incident perturbation. Later on, the heating becomes less enhanced or even reduced because the phase of the reflection has shifted.

[26] Another point that warrants discussion is that a quasisteady state of the ionosphere/thermosphere is not reached in a couple of Alfvén travel times (as might be expected for a collisionless plasma) but takes a much longer time to be established, because of ion-neutral collisions. This can be understood by noting that in the quasi-steady state the Lorentz force from the magnetic perturbation must be balanced by the collisional force from the relative flow. Since the ratio magnetic/velocity perturbation in a single propagating wave is in general quite different from the ratio that corresponds to force balance, several bounces of the perturbations between the magnetopause and the bottom of the ionosphere are required until, by suitable superposition of incident and reflected waves, force balance has been established at each height.

[27] We also consider long-time variations. After a time of about  $1/\nu_{ni}$ , the neutrals have gained sufficient velocity to reduce significantly the plasma-neutral velocity difference, and therefore the heating rate decreases. Figure 7 displays the heating rate as a function of time and altitude for an extended time period. Since the time interval is now long compared to the Alfvén travel time, reflection effects are effectively averaged over, and we plot the time in real (physical) units. Given that the neutral-ion collision frequency (see Figure 1) is maximum in the F-layer, the neutral acceleration takes the shortest time there; accordingly, the heating rate in the F-layer becomes insignificant after about 40 min. In the E layer, on the other hand, where  $1/\nu_{ni}$  is in the order of  $10^6$  s, the heating rate displays no observable decrease within the time period shown.

## 4. Discussion and Conclusions

[28] The heating rates presented in section 3 were obtained from self-consistent numerical solutions of the time-dependent equations for the magnetosphere-ionosphere/thermosphere coupling. Although applied in this initial study only to a highly simplified and artificially limited model, they reveal at least qualitatively some important aspects of the dynamic response of the ionosphere/thermosphere to perturbations from the magnetosphere and/or the solar wind. Compared to the conventional models based on an assumed electric field in the magnetosphere and an ionospheric Ohm's law (heightintegrated, in most cases), the present method of solving the time evolution equations and in particular preserving the time derivative of the perturbation magnetic field (instead of assuming a curl-free electric field), ensures the inclusion of significant wave-propagation and transient effects. Compared to previous wave-based studies [e.g., Hughes and Southwood, 1976; Lysak, 2004], the present work not only resolves the multiple reflections of the waves in the ionosphere but also includes the plasma and neutral dynamics.

[29] The driving physical agent of the ionosphere/ thermosphere heating in our simulation is an imposed convection flow at the top boundary, rather than an imposed electric field or a Birkeland (magnetic field-aligned) current as in conventional treatments. The Birkeland current here is an end effect of the plasma flow and magnetic field perturbations, and the electric field is a result of the plasma motion. It has been shown by *Buneman* [1992] and *Vasyliūnas* [2001] theoretically and by *Tu et al.* [2008] numerically that, as long as the Alfvén speed is small compared to the speed of light, it is the plasma flow that produces the electric field but not the electric field that drives the plasma flow. (The question, does the plasma flow produce the electric field or does the electric field drive the plasma flow? together with related issues has recently been reviewed by *Vasyliūnas* [2011].)

[30] The detailed altitude profiles obtained in the simulation confirm the result of *Vasyliūnas and Song* [2005] that the conventional Joule heating (referred to the neutralatmosphere frame of reference) is more accurately described as frictional rather than Ohmic heating. We find that both the true Joule heating (referred to the plasma frame of reference) and the frictional heating contribute to the ionospheric/thermospheric heating, but the frictional heating is dominant, the Joule heating rate being negligibly small except at the lowest altitudes of the ionosphere and contributing only a tiny fraction of the height-integrated value. The consequent fact that the effective heating rate is simply proportional to the square of the relative velocity between the plasma and the neutrals (and determined primarily by the plasma velocity except at very late times) provides a method of evaluating the heating rate that is physically clearer and mathematically more direct (in comparison to the conventional approach involving ionospheric conductance, electric field, and Birkeland currents).

[31] The heating rate per unit volume is largest in the E-layer of the ionosphere, with a secondary peak in the F-layer. The heating rate per unit mass is concentrated in the F-layer, around 350 km in the present case. This indicates that the heating is more effective in the F-layer compared to that in the E-layer, even though the energy density deposited in the E-layer is larger. Since the F-layer is rich in  $O^+$ , the concentrated heating rate per unit mass indicates a possible mechanism of driving  $O^+$  ion upflow and may provide explanation for the enhanced heavy ion (e.g.,  $O^+$ ) upflow/outflow in response to changes of the solar wind/ IMF when the magnetosphere-ionosphere/thermosphere is in the dynamic stages and the heating can be much stronger than that in the quasi-steady state.

[32] The heating rate approaches a quasi-steady state value about 25 Alfvén travel times after the onset of enhanced flow at the top boundary. This transient interval lasts about 25 s in the present model with the top boundary artificially set at 1000 km; with the top boundary at a realistically higher altitude, it is proportionally longer [Song et al., 2009]. The timescale for the magnetosphere-ionosphere/thermosphere system to reach a quasi-steady state is estimated to be of the order of 10–20 min. The transient stage of the heating is thus comparable to or longer than the timescale of many ionospheric/thermospheric phenomena of interest. The heating rate calculated from the conventional Joule heating expression for a quasi-steady state is a very poor approximation during the first 10 Alfvén travel times and is not really adequate until after about 25 Alfvén travel times. The true heating rate during the first 10 Alfvén travel times is significantly greater than the quasi-steady state value because of the large variations and overshoot of the plasma velocity, due to the wave reflections; the height integrated heating rate can reach values twice as large as the quasi-steady state value. The heating rate thus can be significantly underestimated under the quasi-steady state assumption, particularly if the externally imposed flow is continually varying; this means also that assumption of a potential electric field in the ionosphere may not be valid during the prolonged transient period. Enhanced heating during the dynamic stage must be included to correctly account for the overall effects of the magnetosphere-ionosphere/thermosphere coupling on the ionosphere-thermosphere system.

[33] In our simulation the heating rate undergoes several oscillations during the approach to a quasi-steady state. This resembles the multiple oscillations in the energy flux before reaching the steady state found by *Lysak and Dum* [1983]. The present study, however, focuses on the dynamics and heating within the ionosphere/thermosphere rather than the magnetospheric aspects of magnetosphere-ionosphere coupling as in the work of *Lysak and Dum* [1983] and others. Effects of different types of driving sources and scale-dependence have not been considered and remain a topic for our future studies.

[34] Acknowledgments. This study was supported by the NSF grants AGS-0903777 and ATM-0902965 to the University of Massachusetts

Lowell. The authors thank X. Y. Zhou for his help to the numerical simulations. [35] Robert Lysak thanks the reviewers for their assistance in evaluating this paper.

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P. Song and J. Tu, Center for Atmospheric Research, University of Massachusetts Lowell, Lowell, MA 01854, USA. (paul\_song@uml.edu; jiannan tu@uml.edu)

V. M. Vasyliūnas, Max-Planck-Institut für Sonnensystemforschung, Katlenburg-Lindau, D-37191, Germany. (vasyliunas@mps.mpg.de)