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Anisotropic fluid modeling of ionospheric upflow: Effects of low-altitude anisotropy and thermospheric winds

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Abstract A new anisotropic fluid model is developed to describe ionospheric upflow responses to magnetospheric forcing by electric fields and broadband ELF waves at altitudes of 90–2500 km. This model is based on a bi-Maxwellian ion distribution and solves time-dependent, nonlinear equations of conservation of mass, momentum, parallel energy, and perpendicular energy for six ion species important to $E$, $F$, and topside ionospheric regions. It includes chemical and collisional interactions with the neutral atmosphere, photoionization, and electron impact ionization. This model is used to examine differences between isotropic and anisotropic descriptions of ionospheric upflow driven by DC electric fields, possible effects of low-altitude (<500 km) wave heating, and impacts of neutral winds on ion upflow. Results indicate that isotropic models may overestimate field-aligned ion velocity responses by as much as ~48%. Simulations also show significant anisotropic responses at low altitudes to wave heating for very large power spectral densities, but ion temperature anisotropies below the $F$ region peak are dominated by frictional heating from DC electric fields. Neutral winds are shown to play an important role regulating ion upflow. Thermospheric winds can enhance or suppress upward fluxes driven by DC and BBELF fields by 10–20% for the cases examined. The time history of the neutral winds also affects the amount of ionization transported to higher altitudes by DC electric fields.

1. Introduction

Ion production, loss, and transport in the high-latitude $F$ region ionosphere are regulated by electric fields and auroral precipitation, both of which can lead to strong thermal plasma upflow. Frictional heating-driven upflow (type 1 of Wahlund et al. [1992]) events are typically associated with elevated ion temperatures, strong convection electric fields, and minimal auroral precipitation. In these events, strong convection of the ionosphere through the neutral atmosphere leads to frictional heating of the ions in the $E$ and $F$ regions of the ionosphere, resulting in anisotropic ion distributions [St-Maurice and Schunk, 1979] and large pressure gradients that accelerate ions upward along the field lines [Foster et al., 1998; Zettergren and Semeter, 2012]. Observations of frictional heating-driven upflows often show a lifted $F$ region peak location, low electron densities below 300 km, and modest increases in electron temperature [Wahlund et al., 1992]. In contrast, electron heating-driven upflow (type 2 of Wahlund et al. [1992]) events are associated with auroral precipitation that increases electron densities and temperatures, hence pressure, throughout the $F$ region and topside ionosphere. The electron pressure increase results in a stronger ambipolar electric field which enhances the upward field-aligned flow of plasma. Electron heating-driven upflows are found above auroral arcs, seem to occur more often, and are usually stronger than frictional heating-driven upflows [Foster and Lester, 1996; Wahlund et al., 1992; Ogawa et al., 2003]. Thermal ion upflow mechanisms may not be strong enough to accelerate ions to escape velocities but instead are thought to provide source populations for higher-altitude energization processes. Once ions have been lifted to high altitudes, transverse ion acceleration by broadband ELF waves may give the upflowing ions sufficient energy, which can be converted into parallel momentum through the mirror force, to outflow into the magnetosphere [Kintner et al., 1996; Andre et al., 1998; Moore et al., 1999]. The existence of a multistep process resulting in ionospheric outflow is supported by observations of concurrent ion upflow and outflow drivers (fields, precipitation, ELF waves, etc.) [Yoshida et al., 1999; Lynch et al., 2007; Ogawa et al., 2008; Strangeway et al., 2005]. A variety of modeling studies have established many of the general characteristics of electron heating-driven upflows and outflows driven by auroral processes. The Dynamic Fluid-Kinetic model (DyFK) [Wu et al., 1999] is
a one-dimensional ionospheric model consisting of the Field Line Interhemispheric Plasma model of the ionosphere [Richards and Torr, 1996] coupled with a generalized semikinetic model [Wilson et al., 1990] for higher altitudes where distributions become non-Maxwellian. Studies using this model have shown that soft electron precipitation and perpendicular ion heating can act together to produce intense plasma upflows and outflows [Wu et al., 1999, 2002; Zeng and Horwitz, 2007]. Sadler et al. [2012] have used a three-fluid model (consisting of ion, electron, and neutral fluids) to describe electromagnetic interactions between the ionosphere and the magnetosphere and resulting ion outflow. They found altitude-dependent neutral density enhancements accompany ion upflow driven by soft electron precipitation. Sydorenko and Rankin [2013] developed a multi-fluid model which was used to demonstrate that ion upflow is strongly affected by enhanced ambipolar electric fields produced by soft electron precipitation. Soft electron precipitation is ultimately more efficient at driving ion upflow than high-energy precipitation. High-energy particles penetrate to lower altitudes where heated electrons cool rapidly due to ion-neutral collisions, hence minimizing the upflow response [Su et al., 1999]. Varney et al. [2015] have used an eight-moment fluid model, the Ionosphere/Polar Wind Model, which describes suprathermal ion conic distributions using a separate fluid, to elucidate the details of cusp ion outflow in the presence of wave-particle interactions and significant plasma convection. Changes in ion upflow upstream of the heating region can affect the resulting ion outflow due to convection into and out of the wave heating region.

Frictional heating-driven upflows typically occur less often than electron heating-driven ion upflows [Ogawa et al., 2003]. Zettergren and Semeter [2012] have used a five-moment isotropic fluid ionospheric model to look at contributions of frictional heating, current closure, and ion upflow to auroral density depletions. Their simulations have shown that F region density depletions can form in a matter of minutes due to frictional heating, generation of molecular ions, and enhanced recombination. This depletion process does not significantly inhibit upward fluxes of ions generated by the heating since the upflows are initiated, in the topside ionosphere, on similar timescales. However, any subsequent heating processes occurring in the same region would have access to a more tenuous plasma, which would ultimately impact any upflows that would be generated. Zettergren et al. [2014] and Fernandes et al. [2016] used data from the MICA sounding rocket to drive the ionospheric model developed in Zettergren and Semeter [2012] to examine fine-scale ion upflow and downflow patterns near a series of expansion phase auroral arcs. Observed correlations between ion temperature and the ELF wave power near one of the auroral arcs that the sounding rocket traversed suggested that wave-particle interactions may contribute to temperature anisotropies at altitudes deep in the ionosphere (<400 km altitude). These data also further suggest that wave-particle interactions may provide a significant amount of energy to the ions in the collisional region below 500 km. This observation may be corroborated by a recent study in which the SWARM satellites observed apparent ion temperature anisotropies up to 5 at altitudes as low as 500 km. The anisotropy values far exceed those predicted by theories of DC electric field-driven heating [Archer et al., 2015]. Recent PFISR and RISR experiments show strong evidence of anisotropy in regions of ion frictional heating [Zettergren et al., 2014; Perry et al., 2015]. Significant anisotropies are predicted to occur anytime the DC electric fields exceeds ~50–75 mV/m [St-Maurice and Schunk, 1979; Raman et al., 1981].

While most ionospheric models use an isotropic, collisional fluid description, several theoretical and numerical fluid treatments have included temperature anisotropies. Most of these are 13-moment descriptions [e.g., Zettergren et al., 2010] (based on an expansion about a Maxwellian distribution) or 16-moment descriptions [e.g., Marchaudon and Blelly, 2015] (based on an expansion about a bi-Maxwellian distribution). Both of these formulations include separate parallel and perpendicular pressures and appropriate collision terms; however, only systems based off a bi-Maxwellian distribution are appropriate for large temperature anisotropies [Barakat and Schunk, 1982a]. Transport equations based on a bi-Maxwellian distribution function were first derived by Chew et al. [1956] for a fully ionized, collisionless, anisotropic plasma. Heat flow contributions to the transport equations were neglected and the resulting parallel and perpendicular energy equations were termed “double-adiabatic” energy equations. That work was expanded by Chodura and Pohl [1971] to include both collisionless and collisional transport effects for a fully ionized plasma thus allowing for heat flow, viscosity, and Coulomb (ion-ion) collisions. Demars and Schunk [1979] extended those bi-Maxwellian transport equations to an anisotropic plasma of an arbitrary degree of ionization. Like Chodura and Pohl [1971], Demars and Schunk [1979] also include heat flow and viscosity, but their collision terms were calculated for an arbitrary inverse-power interaction potential which encompassed Coulomb, Maxwell, hard sphere, and resonant charge exchange interactions as special cases. Barakat and Schunk [1982a] extended the transport equations of Demars and Schunk [1979] to include additional physical parameters, such as...
mirror effects, by removing the assumption of straight magnetic field lines. Finally, Blelly and Schunk (1993) analyzed the differences between a 5-, 8-, 13-, and 16-moment transport descriptions by numerically modeling the response of each system in transient situations.

Although a few models resolving anisotropy exist, none has yet been developed specifically to study ion upflow, even though frictional heating-driven upflows, for example, are generated by strong DC fields which will also cause significant anisotropy. This paper presents a new ionospheric model based on a bi-Maxwellian distribution that functions in two spatial dimensions, incorporates all of the ionospheric chemistry and collisional terms needed to properly simulate low-altitude dynamics, and includes possible effects of low-altitude wave-particle interactions. The model accepts, as inputs, the main drivers of ion upflow and outflow: particle precipitation, electric fields, ELF wave power, and neutral winds and densities. This model is critically compared against a Maxwellian model and used in various configurations to assess the possible impacts of low-altitude wave heating and neutral winds on DC electric field driven ion upflows.

2. Anisotropic Fluid Model

Our new anisotropic ionospheric model has been constructed based on a modified 16-moment transport description and is an extension of the isotropic, 5-moment model presented in Zettergren and Semeter [2012] and Zettergren and Snively [2015]. The new model, hereafter referred to as GEMINI-TIA, solves the time-dependent, nonlinear equations of conservation of mass, momentum, parallel energy, and perpendicular energy for six ion species important to the E, F, and topside ionospheric regions: O⁺, NO⁺, N₂⁺, O₂⁺, N⁺, and H⁺. Electrons have also been included using an isotropic description. This model includes chemical and collisional interactions with the neutral atmosphere, as well as the effects of photoionization [Schunk and Qian, 2005] and electron impact ionization [Fang et al., 2008]. Neutral densities and temperatures needed for these calculations are taken from the NRL-MSISE-00 empirical model [Picone et al., 2002] for this study but can also be constrained by models of neutral dynamics (e.g., as in Zettergren and Snively [2013]).

It should be noted that GEMINI-TIA does not solve the full 16-moment set of equations, a task which presents fundamental problems [Palmadesso et al., 1988]. Instead, the system is altered by adopting "equations of state" for the parallel and perpendicular heat fluxes and stresses (viz., separate transport equations for these quantities are not solved). The processes of creating a generalized system of transport equations involves truncating an expansion of the distribution function [Schunk and Nagy, 2000]. This has the side effect of introducing nonphysical wave modes through incomplete phase mixing [Palmadesso et al., 1988; Gombosi and Rasmussen, 1991; Ho et al., 1993]. With the complete phase mixing inherently present in kinetic modeling, a spatial perturbation in one moment damps as the perturbation cascades to ever higher moments and is finally diffused away. Truncating the distribution expansion (and system of equations) stops this cascade and creates nonphysical wave modes in those final, higher-order moments. These waves may then interact with physically meaningful waves with the possibility of creating spurious instabilities. As a relevant example, in the general 16-moment description, the full heat flux moment can become unstable when it has a magnitude on meaningful waves with the possibility of creating spurious instabilities. As a relevant example, in the general 16-moment description, the full heat flux moment can become unstable when it has a magnitude on the order of the thermal energy multiplied by the thermal velocity [Palmadesso et al., 1988]. The equations of state adopted in our model, described below following the model equations, avoid this undesirable situation (Ho et al. [1993], cf. for a similar theoretical treatment).

The transport equations solved in the model are the continuity, parallel momentum, parallel energy, and perpendicular energy equations:

\[
\frac{\partial \rho_i}{\partial t} + \nabla \cdot (\rho_i \mathbf{v}_i) = m_i \dot{P}_i - L_i \rho_i \tag{1}
\]

\[
\frac{\partial (\rho_i \mathbf{v}_i)}{\partial t} - \nabla \cdot \left( (\rho_i \mathbf{v}_i \cdot \mathbf{e}_\parallel) \cdot \mathbf{e}_\parallel \right) = (\rho_i \mathbf{g}) \cdot \mathbf{e}_\parallel - \nabla p_{s,\parallel} \cdot \mathbf{e}_\parallel + (n_i q_i \mathbf{E}) \cdot \mathbf{e}_\parallel
\]

\[
- (p_{s,\parallel} - p_{s,\perp}) \nabla \cdot \mathbf{e}_\parallel - \sum_j \frac{3 v_j}{4 \pi k_B} \left[ 2 k_B n_j m_j \frac{\sigma_{s,j,\parallel,002}}{\sigma_{s,j,\parallel}} (\mathbf{v}_j - \mathbf{v}_i) \right]
\]

\[
- \sum_n n_i m_i v_{in} (\mathbf{v}_i - \mathbf{v}_n) \tag{2}
\]
The ion stress tensor (the term “stress” is used, in this paper, in the sense defined in the Boltzmann collisional theory) has been added to include chemical loss processes. Equations of state are adopted for the heat fluxes to the grid. In Equations (1)–(4) are based off of Barakat and Schunk [1982a, equations (11)–(14) and (29)] with a few modifications. In equation (1), the continuity equation, the species production rate, $P_s$, has been added to include ion creation from chemical production, photoionization, and impact ionization. The loss frequency term, $Q_l$, has been added to account for chemical loss processes. Equations of state are adopted for the heat fluxes to prevent our fluid system from needing additional transport equations describing the time-dependent evolution of heat fluxes. For the region of the model where collisions are significant (both Maxwell and Coulomb), traditional Fourier’s law descriptions are used.

$$\frac{\partial p_{s\parallel}}{\partial t} + \nabla \cdot (p_{s\parallel} \mathbf{v}_s) = -2p_{s\parallel} (\nabla \cdot \mathbf{v}_s) - \nabla \cdot (h_{s\parallel} \mathbf{e}_\parallel) + 2h_{s\perp} (\nabla \cdot \mathbf{e}_\parallel) + m_0 k_b \sum_j \left[ \frac{\sigma_{s\perp}}{\sigma_{s\parallel}} \left( T_{s\parallel} - T_{s\perp} \right) \right]$$

$$+ m_j \left( \frac{2\pi}{3} (\mathbf{v}_j - \mathbf{v}_s)^2 + 2\sigma_{s\perp} (I_{000} - I_{002}) \right)$$

$$- \sum_n \frac{n_m v_{\text{ms}}}{(m_s + m_n)} \left[ 2k_b (T_{s\parallel} - T_{s\perp}) - 2m_n (\mathbf{v}_s - \mathbf{v}_n)_\parallel \right]$$

$$+ m_n Q_{in}^{(2)} \left[ 2k_b (\sigma_{s\perp} - \sigma_{s\parallel}) - (\mathbf{v}_s - \mathbf{v}_n)^2 + 3(\mathbf{v}_s - \mathbf{v}_n)_\parallel^2 \right]$$

(3)

$$\frac{\partial p_{s\perp}}{\partial t} + \nabla \cdot (p_{s\perp} \mathbf{v}_s) = -p_{s\perp} (\nabla \cdot \mathbf{v}_s) - \nabla \cdot (h_{s\perp} \mathbf{e}_\perp) + W_{s\perp} - h_{s\parallel} (\nabla \cdot \mathbf{e}_\parallel) + m_0 k_b \sum_j \left[ \frac{\sigma_{s\parallel}}{\sigma_{s\perp}} \left( T_{s\parallel} - T_{s\perp} \right) \right]$$

$$+ m_j \left( \frac{4\pi}{3} (\mathbf{v}_j - \mathbf{v}_s)^2 + 2\sigma_{s\parallel} (I_{002} - I_{020}) \right)$$

$$- \sum_n \frac{n_m v_{\text{ms}}}{(m_s + m_n)} \left[ 2k_b (T_{s\parallel} - T_{s\perp}) - m_n (\mathbf{v}_s - \mathbf{v}_n)^2 \right]$$

$$+ m_n Q_{in}^{(2)} \left[ 2k_b (\sigma_{s\parallel} - \sigma_{s\perp}) - 2(\mathbf{v}_s - \mathbf{v}_n)^2 + 3(\mathbf{v}_s - \mathbf{v}_n)_\parallel^2 \right]$$

(4)

In these equations $p_s$ is the mass density, $n_s$ is the number density, $m_s$ is the mass, and $v_s$ is the drift velocity of species $s$. Parallel and perpendicular pressures are $p_{s\parallel} = p_s k_b T_{s\parallel}/m_s$ and $p_{s\perp} = p_s k_b T_{s\perp}/m_s$, respectively, with $k_b$ as the Boltzmann constant, $T_{s\parallel}$ is the parallel temperature, and $T_{s\perp}$ is the perpendicular temperature. The charge of each species is represented by $q_s$. $E$ is the electric field, $g$ is the gravitational field, and $\mathbf{e}_\parallel$ is a unit vector along the geomagnetic field. Additional variables (i.e., collision term quantities) within these transport equations are defined in greater detail in paragraphs that follow.

Equations (1)–(4) are based off of Barakat and Schunk [1982a, equations (11)–(14) and (29)] with a few modifications. In equation (1), the continuity equation, the species production rate, $P_s$, has been added to include ion creation from chemical production, photoionization, and impact ionization. The loss frequency term, $Q_l$, has been added to account for chemical loss processes. Equations of state are adopted for the heat fluxes to prevent our fluid system from needing additional transport equations describing the time-dependent evolution of heat fluxes. For the region of the model where collisions are significant (both Maxwell and Coulomb), traditional Fourier’s law descriptions are used.

$$h_{s\parallel} = -\frac{2}{3} \lambda_s V_{s\parallel} T_{s\parallel} \cdot \mathbf{e}_\parallel$$

(5)

$$h_{s\perp} = -\frac{2}{3} \lambda_s V_{s\perp} T_{s\perp} \cdot \mathbf{e}_\perp$$

(6)

$$h_e = -\frac{2}{3} \lambda_e V_{e\parallel} T_e \cdot \mathbf{e}_\parallel - \beta_e \mathbf{J} \cdot \mathbf{e}_\parallel$$

(7)

where $\lambda_s$ is the thermal conductivity for ion species $s$, which is calculated using the parallel temperature [e.g., Singh, 1992], $\lambda_e$ is the electron thermal conductivity, $\beta_e$ is the thermo-electric coefficient, and $\mathbf{J}$ is the current density. It is assumed that only the parallel component of the heat flux exists (viz., only the parallel transport of parallel and perpendicular thermal energy is considered). This collisional heat flux is smoothly tapered to zero for the upper altitudes of our model grid where the Knudsen number is greater than one (indicated to a transition to a collisionless plasma). This transition occurs at roughly 2500 km altitude for the conditions examined in this study, so the higher-altitude portions of the grid mostly serve as a buffer zone to prevent boundary conditions from affecting the numerical solutions. Finally, the new model neglects the ion stress tensor (the term “stress” is used, in this paper, in the sense defined in Barakat and Schunk [1982a]).
Most ionospheric models do not include ion stress since it significantly complicates application of the transport equations (see also the discussion above regarding the full 16-moment system of equations). Additionally, simpler descriptions of ion stress like the Navier-Stokes approximations are not likely to be valid at altitudes where ion stress matters [cf., Schunk, 1975]. Note that while we exclude stress effects, we again emphasize that we do retain separate temperatures for the parallel and perpendicular directions.

The multifluid system of equations (1)–(4) is closed through an electrostatic treatment of the auroral currents:

\[
\nabla_\perp \cdot (\sigma_\perp \cdot \nabla_\perp \Phi) + \nabla_\parallel \cdot (\sigma_\parallel \nabla_\parallel \Phi) = \nabla_\perp \cdot \left( \sum_s n_s m_s \mu_{s,\perp} \cdot \mathbf{v}_s,\perp \right)
\]

where \( \Phi \) is the electric potential corresponding to the resistive part of the electric field (viz., excluding ambipolar field contributions), \( \sigma_\perp \) is the perpendicular conductivity tensor, and \( \sigma_\parallel \) is the parallel conductivity—both defined in Zettergren and Semeter [2012, equations (15) and (16)]. Since the electron mobility is much higher than the ion mobility in the parallel direction, the parallel ion drift contributions to current are neglected.

Many of the assumptions inherent in this electrostatic treatment are discussed in greater detail in Zettergren and Semeter [2012] and Zettergren and Snively [2015]. The electric field in the perpendicular direction is found from \( \mathbf{E} = -\nabla_\perp \Phi \), while the parallel electric field is obtained by superposing the resistive part of the field (as computed from equation (8)) and the ambipolar electric field [see also Zettergren and Semeter, 2012], i.e.,

\[
\mathbf{E}_\parallel = -\nabla_\parallel \Phi + \frac{1}{n_e q_e} \nabla_\parallel \rho_e.
\]

A static geomagnetic field, presently a tilted dipole model, is used when calculating the conductivities and mobilities needed for equations (8), the ion mobilities in equation (10), and the metric coefficients needed to solve the system of equations defined by equations (1)–(8) [Huba et al., 2000].

Consistent with the use of an electrostatic treatment, a steady state momentum approximation is used for the perpendicular ion and electron drifts.

\[
\mathbf{v}_{s,\perp} = \mu_{s,\perp} \cdot \left( \mathbf{E}_\perp + \frac{m_s v_s}{q_s} \mathbf{v}_{s,\perp} \right)
\]

where \( \mu_{s,\perp} \) is the ion mobility tensor and \( v_s \) is the total ion-neutral collision frequency (Zettergren and Semeter [2012], defined in their equations (9) and (11), respectively).

Electrons are treated differently from the ions in our new model [e.g., Zettergren and Snively, 2015]. The electron number density is found via quasi-neutrality

\[
n_e = \sum s_n_s,
\]

and the electron velocity is calculated from the steady state current density

\[
\mathbf{v}_{e,\parallel} = -\frac{1}{n_e q_e} \left( \sum_s n_s q_s \mathbf{v}_{s,\parallel} - \mathbf{J}_\parallel \right).
\]

where \( \mathbf{J}_\parallel = \sigma_0 (-\nabla_\parallel \Phi) \). Unlike the ions, the electrons are considered to be isotropic so only a single transport equation is solved for the electron energy:

\[
\frac{\partial (\rho_e \varepsilon_e)}{\partial t} + \nabla \cdot (\rho_e \varepsilon_e \mathbf{v}_e) = \rho_e \left( \nabla \cdot \mathbf{v}_e \right) + \nabla \cdot \mathbf{h}_e + \frac{Q_e}{(\gamma_e - 1)} - \frac{1}{\gamma_e - 1} \sum_s \rho_e q_s v_s \left[ 2 (\tau_e - \tau_s) - \frac{2m_v}{3k_b} (\mathbf{v}_e - \mathbf{v}_s)^2 \right]
\]

where \( \varepsilon_e = \rho_e / [(\gamma_e - 1) \rho_e] \) is the specific internal energy. This electron energy equation differs from the ion energy equations in that it also includes inelastic cooling terms [Schunk and Nagy, 1978] and heating by
photoelectrons in the \( Q_e \) term [Swartz and Nisbet, 1972]. There is also a thermoelectric component in the heat flux term as shown in equation (7) [e.g., Schunk and Nagy, 1978; Zettergren and Snively, 2015].

Some assumptions encoded in the model’s system of equations deserve further comment. As discussed, the electrostatic equation (equation (8)) assumes a steady state momentum balance for the principle charge carriers (electrons in the parallel direction and all type of particles in the perpendicular direction). The justification for considering only the parallel electron current is based on the fact that the parallel electron mobility is much larger than that of the ions. Because the ion current parallel to the geomagnetic field is neglected, we are able to use a different formulation (i.e., the full time-dependent, nonlinear momentum equation) for the parallel ion momentum. The reason for doing this is that the process of interest to the present study, ion upflow and outflow, requires a time-dependent, nonlinear description to properly capture the complicated behavior of very large ion heating events. Because ion momentum contributes negligibly to electrical current, using a time-dependent formulation does not invalidate the assumptions underlying our current continuity equation. The validity of treating the parallel ion momentum in a time-dependent manner while simultaneous using a steady-state for the parallel electron momentum is based on the fact that in the presence of collisions, electrons will attain a steady state more quickly than the ions. This is a consequence of the fact that the ion-electron collision frequency is much smaller than the electron-ion collision frequency [Schunk and Nagy, 2000, chap. 4].

2.1. Wave Heating

The resonant heating term, \( \dot{W}_{s,\perp} \), in equation (4) parameterizes the acceleration of ions by transverse plasma waves. This gyroresonant (cyclotron) energy transfer requires low ion-neutral collision rates and is therefore rarely observed below 500 km except in extreme cases [e.g., Whalen et al., 1978]. Because this type of heating occurs primarily in collisionless regions, the resulting ion distributions remain highly anisotropic and are accelerated by the mirror force, attaining large field-aligned velocities high above the heating region. The present form of the model uses an empirical specification of this heating term:

\[
\dot{W}_{s,\perp}(\omega) = 2 \rho_s \left( \frac{\eta q_s^2}{4 m_s^2} \right) |E_0|^2 \left( \frac{\omega}{\omega_o} \right)^{-\alpha} \tag{14}
\]

where \( \omega \) is the local gyrofrequency for each ion, \( \eta \) is the fraction of the wave field which is left-hand polarized, assumed to be 0.125 [Chang et al., 1986], \( \alpha \) is the spectral power index, assumed to be 1.7 [Crew et al., 1990], and \( |E_0|^2 \) is the wave power spectral density at some reference frequency \( \omega_o \), where \( \omega_o \) is taken to be 6.5 Hz [Zeng and Horwitz, 2008; Zeng et al., 2006; Retterer et al., 1983]. The wave heating term is evaluated at the local gyrofrequency of each ion species for every point in the simulation and smoothly tapered to zero at the top of the simulation, similar to what was done in Wu et al. [1999], to prevent boundary condition artifacts. Note that even though the wave power spectral density is specified as constant with altitude [Bouhram et al., 2003] the heating rate will still be altitude dependent since the ion gyrofrequency and density change with altitude.

2.2. Coulomb Collisions

The Coulomb collision terms used to describe ion-ion interactions (the summations over index “\( j \)” in equations (2)–(4)), are taken from Biely and Schunk [1993, equations (61), (63), (66)]. These terms were derived for small stress and heat flows, low-speed plasma flows where the species drift velocity differences are small in comparison to the thermal speeds, arbitrary difference between species temperatures, and an arbitrary difference between parallel and perpendicular temperatures for the same species. This summation represents the response to a collisional interaction between ion species \( s \) and \( j \). Within the Coulomb collision summations the following definitions are used:

\[
\sigma_{s||} = \frac{T_{s||}}{m_s} + \frac{T_{j||}}{m_j} \tag{15}
\]

\[
\sigma_{s\perp} = \frac{T_{s\perp}}{m_s} + \frac{T_{j\perp}}{m_j}. \tag{16}
\]

The collision frequency, \( \nu_{sj} \), for Coulomb collisions is

\[
\nu_{sj} = B_j n_j \sigma_{s||}^{-1/2} \sigma_{s\perp}^{-1} \tag{17}
\]
where $B_{ij}$ is tabulated in Schunk and Nagy [2000, Table 4.3] for species relevant to this work. The $I_{LMN}$ quantities are defined in Demars and Schunk [1979] as

$$I_{LMN} = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \frac{x_1^L x_2^M x_3^N}{(x_1^2 + x_2^2 + x_3^2)^{3/2}} e^{-\left(x_1^2 + x_2^2 + x_3^2\right)^{3/2}} \, dx_1 \, dx_2 \, dx_3$$

which are the same as the $K_{LMN}$ quantities in Chodura and Pohl [1971] multiplied by $\pi$. This triple integral can be converted, as shown in Demars and Schunk [1979], to the relationships listed in Chodura and Pohl [1971, Table 1]. From that table, the $I_{LMN}$ relationships useful for this study are reproduced, for easy reference, in Table 1 of this paper.

As indicated by the Coulomb collision summation in equation (2), the rate of change of momentum is proportional to ion-ion drag. Within the Coulomb collision summations in equations (3) and (4), the rate of change of parallel and perpendicular energy is proportional the terms which, left to right, represent heat exchange, frictional heating, and parallel-perpendicular heat transfer within a species having a temperature anisotropy.

### 2.3. Maxwell Interactions

The Maxwell interaction terms describing the ion-neutral interactions (summations over index “n” in equations (2)–(4)) are taken from Barakat and Schunk [1982a, equations (33)–(35)] and Demars and Schunk [1979, equations (23a)–(23c)]. These terms were derived for arbitrary relative drifts between different species, arbitrary temperature anisotropies for each species, and arbitrary temperature differences between species. Each summation represents the response from a collisional interaction between an ion of species $s$ and a neutral of species $n$. The neutral species currently included are O, N, O$_2$, and H. Within the Maxwell interaction terms in the following definitions are used:

$$\sigma_{sn,\parallel} = \frac{T_{s,\parallel}}{m_s} + \frac{T_{n,\parallel}}{m_n}$$

$$\sigma_{sn,\perp} = \frac{T_{s,\perp}}{m_s} + \frac{T_{n,\perp}}{m_n}$$

The collision frequency, $\nu_{sn}$, suitable for Maxwell interactions is

$$\nu_{sn} = \frac{C_{sn} n_s \sigma_{sn,\parallel}^{-1/2} \sigma_{sn,\perp}^{1/2}}{\sigma_{sn,\parallel} \sigma_{sn,\perp}}$$

where $C_{sn}$ is tabulated in Schunk and Nagy [2000, Table 4.4] for nonresonant interactions and in Schunk and Nagy [2000, Table 4.5] for resonant interactions.

The $Q_{sn}^{(2)}/Q_{sn}^{(1)}$ ratios used in equations (3) and (4) have been updated from the original values used by Barakat and Schunk [1982b] to take into consideration the relative energy between colliding particles. As the relative energy of colliding particles changes, the nature of their interaction varies as well. At low energies, collisions are dominated by the polarization attraction between ions and neutrals. At higher energies, a repulsive $1/r^{12}$ potential takes over and changes the generalized collision cross section. This variation in the collision cross section was explored in Gaimard et al. [1998] through their comparison between analytical results and Monte Carlo simulations. In analytical calculations the speed dependence of the collision cross section has to be neglected. This restriction is not present in the Monte Carlo simulations where the cross section is allowed to change with the relative energy of the colliding particles producing more accurate temperatures. Gaimard et al. [1998] developed a set of cross sections that vary with respect to the DC electric field strength in such a way as to take account of the modifications in the nature of the ion-neutral interactions as the relative energy of the colliding particles changes. A $Q_{sn}^{(2)}/Q_{sn}^{(1)}$ ratio of 0.85 is a good average for almost all of the ion-neutral interactions.

| Table 1. The $I_{LMN}$ Quantities for Coulomb Collisions$^a$ |
|-------------|----------|----------|----------|
| $L$ | $M$ | $N$ | $I_{LMN}$ |
| 2 | 0 | 0 | $(-1 + (1 + \lambda)|\text{atan}(\sqrt{\lambda})/\sqrt{\lambda}|)/\lambda$ |
| 0 | 0 | 2 | $2(1 - |\text{atan}(\sqrt{\lambda})/\sqrt{\lambda}|)/\lambda$ |

$^a \lambda = (\sigma_{s,\perp}/\sigma_{s,\parallel}) - 1.$
interactions [Gaimard et al., 1998], but it should be noted that the charge exchange cross section is much larger, compared to the polarization cross section for ion-neutral collisions. The $Q_{\text{ex}}^{(2)} / Q_{\text{ex}}^{(1)}$ ratio value for $\text{O}^+\text{--O}$ resonant interactions will therefore have a strong dependence on the electric field. A simple polynomial fit to the values from Gaimard et al. [1998, Table 1] for various electric fields is used in the model for the $\text{O}^+\text{--O}$ resonant interaction. This polynomial fit given by the following:

$$\frac{Q_{\text{ex}}^{(2)}}{Q_{\text{ex}}^{(1)}} = 6.284 \times 10^{-6} E^2 - 2.833 \times 10^{-3} E + 0.5348. \quad (22)$$

where $E$ is the magnitude of the perpendicular electric field.

As indicated by the Maxwell interaction summation in equations (2), the rate of change of momentum from this process is proportional to ion-neutral drag. Within the Maxwell interaction summations in equations (3) and (4), the rate of change of parallel and perpendicular energy is proportional the terms which, left to right, represent heat transfer between different species, frictional heating, heat transfer between parallel and perpendicular directions and two more terms of frictional heating.

### 2.4. Numerical Methods

Equations (1)–(4), (10) (ion dynamics), (11)–(13) (electron dynamics), and (8) (electric potential), are solved to define the ionospheric dynamics within this model. A split time step procedure is used to separate the advection portion of the fluid equations (left-hand side of equations (1)–(4) and (13)) from the diffusion and source/loss terms (right-hand side). The advection component is solved using a flux-limited finite volume method (MC flux limiter). The remaining terms in the ion continuity and momentum equations are source/loss terms which are solved using an exponential time differencing method (ETD). The remaining terms in both the ion and electron energy equations (equations (3), (4), and (13)) contain diffusion and source/loss terms. Within the electron energy equation, the diffusion parts are solved using a trapezoidal backward difference method (TRBDF2) and the source/loss parts are solved using a Runge-Kutta (RK) method for the compression term and ETD for the remaining terms, including collisions. For the ions, the compression term in both parallel and perpendicular energy equation is also solved using a RK method. The diffusion and source/loss terms (collision terms) for both the parallel and perpendicular direction are solved simultaneously, rather than in separate substeps, due to the strong coupling between parallel and perpendicular heat fluxes, frictional heating, and heat exchange. The solution for the ion diffusion and source substep is done with a first-order backward difference in time and second order centered difference in space (the backward time centered space method). The electric potential, equation (8), uses a finite difference technique to generate a sparse system of linear equations that is solved through LU factorization [Davis, 2004]. The perpendicular ion velocity, electron density, and electron velocity from equations (10)–(12) are solved algebraically once the relevant variables have been calculated.

### 3. Results

A similar model configuration is used for each simulation conducted as part of this study. Each simulation was run for 10 min with a two second output cadence. The adaptive time stepping of the model is such that stability is retained at every time step, and a typical time step is $\sim 0.5$ s. The model was run to a steady state for the initial conditions, and all of the simulations start at 15 UT creating a consistent set of background conditions for each run. The simulation results presented here use a dipole mesh [Huba et al., 2000; Zettergren and Snively, 2015] spanning L shells 12–16, centered roughly on the location of the Sondrestrom research facility on the west coast of Greenland ($\sim 67^\circ$, $309^\circ$), a location of interest for ion upflow [Semeter et al., 2003; Zettergren et al., 2008; Sánchez and Stømme, 2014]. For this study, GEMINI-TIA was run for multiple combinations of DC electric fields, transverse wave heating, and neutral winds implemented using the configurations described in paragraphs below.

The DC electric field is applied using a Gaussian envelope centered in the domain to prevent side boundary condition artifacts. This electric field is imposed using Dirichlet boundary conditions with a topside potential specified as follows:

$$\Phi(x_2) = E_{\text{OL}} \frac{h_2 c \sqrt{x_2}}{2} \text{erf} \left( \frac{x_2 - b}{c} \right) \quad (23)$$
where \( E_{0z} \) is the strength of the DC electric field (V/m), \( c \) controls the width of the region of electric field, here set to 1/7th of the domain, \( b \) is the location of the center field line of the simulation, and \( h_1 \) is the metric factor corresponding to the L shell dimension of the model.

The wave heating term, \( W_{\perp} \), depends on an adjustable wave power spectral density parameter \([E_0]^2\) (V²/m²/Hz) and a reference \( O^+ \) gyrofrequency \( \omega_0 \) used here as 6.5 Hz [Zeng et al., 2006]. Similar to how a Gaussian spatial envelope is used for the DC electric field, the wave heating term is constrained to prevent excessive energization of the ions near boundaries. Perpendicular to the field lines, the standard deviation of this envelope is 1/6th of the domain size in that direction. Parallel to the field line, at 1/20th of the distance (measured in terms of the field-aligned variable \( x_1 \); \( q \) in the notation of Huba et al. [2000]), from the top of the simulation, a hyperbolic tangent is used to quickly, over 1/100th of the range, taper off the heating term to prevent spurious boundary interactions. This transition region is well outside the area of interest to this study so the results are not impacted by the selection of altitude at which the wave heating is removed.

For simulations that include geographic northward (or southward) neutral winds, these are specified using components parallel and perpendicular to the magnetic field. The parallel neutral wind component is given by the following:

\[
v_{n,x_1} = v_{n,0} \cos(i) \left( \frac{1}{2} + \frac{1}{2} \tanh \left( \frac{z - z_1'}{\Delta z_1} \right) \right),
\]

where here \( v_{n,0} \) is the northward geographic neutral wind and \( i \) is the inclination of the magnetic field lines from the horizontal. Thus, \( v_{n,0} \cos(i) \) determines the component of the geographic neutral wind along the magnetic field lines, the \( x_1 \) direction. The hyperbolic tangent is used to, over an altitude span of \( \Delta z_1 = 10 \) km, centered at \( z_1' = 90 \) km, taper the winds to full strength preventing lower boundary artifacts. The perpendicular neutral wind component uses a similar relationship:

\[
v_{n,x_2} = v_{n,0} \sin(i) \left( \frac{1}{2} + \frac{1}{2} \tanh \left( \frac{z - z_1'}{\Delta z_1} \right) \right),
\]

where \( v_{n,0} \sin(i) \) determines the component perpendicular to the magnetic field lines, the \( x_2 \) direction. Geographically eastward (and westward) winds, which here are parallel (and antiparallel) to the \( \mathbf{E} \times \mathbf{B} \) drift direction, do not need to be broken into components and are described by

\[
v_{n,x_3} = v_{n,0} \left( \frac{1}{2} + \frac{1}{2} \tanh \left( \frac{z - z_1'}{\Delta z_1} \right) \right) e^{-\frac{q x_1^2 - y^2}{2 a^2}}.
\]

In this equation, \( v_{n,0} \) is the geographic eastward neutral wind, \( c \) is the standard deviation, set here to 1/7th of the domain, and \( b \) is the location of the center field line.

3.1. Comparison of Maxwellian and Bi-Maxwellian Simulations

GEMINI-TIA, the new model developed as part of this work, and its isotropic, parent version (described most recently in Zettergren and Snively [2015, Appendix A]) were first run with identical drivers and initial conditions to assess and clearly illustrate differences in the ionospheric response to frictional heating under different transport formulations. The models were run with a DC electric field, \( E_{0x} \), of 80 mV/m and no wave heating for this comparison. Frictional heating resulting from this strong electric field leads to anisotropies not resolved in the Maxwellian model [St-Maurice and Schunk, 1979] and will result in differences in the simulated upflows.

The \( O^+ \) flux and velocity parallel to the geomagnetic field from both simulations are shown in Figure 1. Figures 1a–1d show snapshots of the model output after the DC electric field has been applied for 250 s with the isotropic model responses in Figures 1a and 1c and the anisotropic model responses in Figures 1b and 1d. Figures 1a and 1b compare the \( O^+ \) fluxes and Figures 1c and 1d compare field-aligned velocity. In general, a significant difference in the Maxwellian versus bi-Maxwellian response can be noted. The isotropic model \( O^+ \) flux (Figure 1a) is 48% larger at 1000 km (with a value of \( 3.1 \times 10^{13} \) m⁻² s⁻¹) than the anisotropic model's \( O^+ \) response (Figure 1b) of \( 2.1 \times 10^{13} \) m⁻² s⁻¹. The isotropic model \( O^+ \) field-aligned velocity (Figure 1c) is 33% larger at 1000 km (with a value of ∼600 m/s) at this point in time than the anisotropic velocity response (Figure 1d) of ∼450 m/s. Figure 1e, shows how ion fluxes at 1000 km evolve over the duration of the simulation on the center field line of the grid. It takes approximately 4 min for the main ion perturbation to reach...
Figure 1. A comparison of the O\textsuperscript{+} flux and field-aligned velocity between the new 16-moment (anisotropic) model and the parent 5-moment (isotropic) model after 250 s of an applied DC electric field of 80 mV/m. (a) Isotropic O\textsuperscript{+} flux, (b) anisotropic O\textsuperscript{+} flux, (c) isotropic field-aligned velocity, (d) anisotropic field-aligned velocity, and (e) both the isotropic and anisotropic O\textsuperscript{+} flux at 1000 km, on the center field line of the simulation, for the entire duration of the simulation. There is a significantly larger response in the isotropic model compared to the anisotropic model, 48% larger at 250 s at 1000 km. The field-aligned velocity is also larger with a 33% increase in the isotropic velocity response when compared to the anisotropic response. Figure 1e highlights the consistency of the isotropic flux response to be larger than the anisotropic flux response.

In the anisotropic case, heating from the DC electric field at lower altitudes supports that rate of upflow for several more minutes until the ionospheric plasma pressure begins to reestablish a force balance resulting in a decrease of the flux. In the isotropic model, the larger flux rate begins to decrease right after peaking at \(\sim 240 \) s but still remains larger than the anisotropic model fluxes for the remainder of the simulation.

The difference in the ion upflow response in the two models is due to the fact that the frictional heating, in reality (and in the bi-Maxwellian-based model), leads to a larger perpendicular temperature than parallel temperature. The average temperature in the bi-Maxwellian simulation, \( T_s = \frac{1}{3} T_{||} + \frac{2}{3} T_{\perp} \), is very similar to what is simulated by the Maxwellian model, there is just a different partitioning of the energy between the parallel and perpendicular directions. Specifically, the Maxwellian model assumes equal partitioning in both directions, while the anisotropic model correctly accounts for a larger fraction of the energy being distributed into the field-perpendicular direction. Ion upflow in the anisotropic model depends on the parallel pressure...
Figure 2. O\(^+\) responses to a DC electric field of \(E_0 = 80\) mV/m, wave heating of \(|E_0|^2 = 0.3\) (mV/m\(^2\))/Hz, and both energy sources applied simultaneously. Simulation results are plotted as a function of time and altitude for the center geomagnetic field line of the grid. (1a–1c) Electron density, parallel velocity, and temperature anisotropy, respectively, for the simulation with just the DC electric field. (2a–2c) Electron density, parallel velocity, and temperature anisotropy, respectively, for the simulation with just wave heating. (3a–3c) Electron density, parallel velocity, and temperature anisotropy, respectively, for the simulation with a DC electric field and wave heating.

gradient term in equation (2), while it depends on the average pressure gradient term in the isotropic model [Zettergren and Snively, 2015, equation (A8)]. Hence, the difference in flux responses in these models is due to the overestimation of the pressure gradient force by the isotropic model. It is worth noting that the anisotropic model contains a mirror force term, which is several orders of magnitude smaller than the pressure gradient term at the altitudes considered in this study and there are also a few differences present in the way the collisions are described; however, the difference in the pressure gradient terms in the bulk momentum transport equations is the primary cause of the smaller upward field-aligned velocities and weaker O\(^+\) flux response in the anisotropic model versus the isotropic model. This exercise shows that isotropic fluid models may overestimate ionospheric velocities and the amount of plasma supplied to higher altitudes by as much as 48% at 1000 km (the case shown here) and illustrates that anisotropies significantly affect the intensity of frictional heating-driven upflow. As a final note, a modest enhancement in ion upflow begins almost as soon as the simulation starts. This is an indirect effect, discussed in detail below, that results from heat transfer from the frictionally heated ions to the electrons at ionospheric altitudes.

3.2. Effects of Low-Altitude Wave Heating on Ion Upflow

GEMINI-TIA is used here to study the impacts of DC electric fields, wave heating, and the synergistic effects of both processes in a set of three simulations. Presented in this section are as follows: (1) a run with just a DC
electric field, (2) a run with just wave heating, and (3) a run with both a DC electric field and wave heating. Figure 2 shows ionospheric state parameters for each of these cases, extracted along the center geomagnetic field line from the simulation, as a function of time and altitude. The first column of plots contain parameters, from a simulation that used $E_0 = 80 \text{ mV/m}$ (constant for the full 10 min duration of the simulation). In descending order these parameters are $O^+$ density, field-aligned velocity, and temperature anisotropy factor (defined as $T_\parallel/T_\perp$). The second column contains the same parameters for a simulation that used $|E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz}$ (constant for the entire simulation). The third column also contains the same parameters taken from a simulation that used both $E_0 = 80 \text{ mV/m}$ and $|E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz}$ (again, constant for the duration of the simulation).

There is a distinct density decrease in the $F$ region in the simulations that include a DC electric field, cf. Figure 2, panels 1a and 3a. This density decrease is due to conversion of $F$ region $O^+$ into molecular ions (a very temperature-sensitive process) [St.-Maurice and Laneville, 1998] which recombine quickly [Schunk, 1975; Zettergren and Semeter, 2012]. The average temperature of the model is used within the reaction coefficients to account for the fact that particles of all pitch angles undergo these reactions. By comparison, the simulation with only wave heating shows an $F$ region density, in panel 2a, that is relatively unchanged with time. The wave heating primarily impacts ion populations above the $F$ region peak and does not greatly affect $N_mF_2$. The limited altitude region of ionospheric response to wave heating is also seen in the field-aligned velocity in Figure 2, panel 2b, which only shows significant responses at the highest altitudes (>1000 km). In this example (with only wave heating) there is a maximum field-aligned velocity of approximately 400 m/s at 2500 km by the end of the simulation. Both panel 1b and 3b have larger field-aligned velocity response (driven by the DC electric field) with the case shown in Figure 2, panel 3b, having the largest velocities which result from the combined effects of frictional heating and wave heating. In this case, the field-aligned velocity at 2500 km at the end of the simulation is 1200 m/s, panel 3b, a 200 m/s increase from the 1000 m/s seen in panel 1b, the case with just the DC electric field-driven frictional heating. The DC electric field-driven anisotropies are primarily at lower altitudes, Figure 2, panel 1c, and the wave heating-driven anisotropies are at higher altitudes, Figure 2, panel 2c. In general, the model shows that the effectiveness of cyclotron wave heating is mitigated, to a degree, by the presence of collisions at the lower altitudes, both ion-ion and ion-neutral. As a result the wave heating-driven temperature anisotropies are only present down to 500 km for this level of wave heating. The simulation that uses both a DC electric field and wave heating has significant temperature anisotropies throughout the entire altitude range of ~150–2500 km.

To further examine the synergistic effects of frictional heating and wave heating at low altitudes, a set of 12 simulations with different combinations of these parameters has been run; Table 2 summarizes these configurations. For purposes of comparison, a control simulation where no drivers are applied (case I listed in Figures 3 and 4 and Table 2) is also included. Figure 3a shows the temperature anisotropy response, and Figure 3b shows the field-aligned velocity response of $O^+$ extracted along the center geomagnetic field line at 30 s for each of the 12 simulations. Note that the parallel and perpendicular energy transport equations naturally create a modest temperature anisotropy at altitudes greater than 600 km even under equilibrium conditions (see the control simulation, case I in Figure 3a). The reference value for the wave power spectral density, $|E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz}$ (case II), creates an increase in the temperature anisotropy seen down to 500 km. This is the nominal minimum altitude that it is normal to see wave heating effects at since the lower altitudes are highly collisional [Archer et al., 2015]. Increasing the wave power spectral density increases the anisotropy factor at high altitudes and also serves to increase the penetration depth of the wave heating effects into the ionosphere. Using $|E_0|^2 = 3.0 \text{ (mV/m)}^2/\text{Hz}$, a strong value within the bounds of the observations, creates significant anisotropies down to 300 km and an extreme value of $|E_0|^2 = 10.0 \text{ (mV/m)}^2/\text{Hz}$, selected for illustrative purposes, can create an observable temperature anisotropy down to 250 km after 30 s of wave heating. At that point in time, the main field-aligned upflow perturbation has reached, on average, 520 km altitude as seen by the location of the peak velocities of Figure 3b. Wave heating also increases upflow

### Table 2. The Wave Power Spectral Density, $|E_0|^2$, and DC Electric Field, $E_0$, Used in Simulations I–XII Plotted in Figures 3 and 4

<table>
<thead>
<tr>
<th>Simulation</th>
<th>I</th>
<th>II</th>
<th>III</th>
<th>IV</th>
<th>V</th>
<th>VI</th>
<th>VII</th>
<th>VIII</th>
<th>IX</th>
<th>XI</th>
<th>XII</th>
</tr>
</thead>
<tbody>
<tr>
<td>$</td>
<td>E_0</td>
<td>^2$ (mV/m)$^2$/Hz</td>
<td>0.0</td>
<td>0.3</td>
<td>3.0</td>
<td>10.0</td>
<td>0.0</td>
<td>0.3</td>
<td>3.0</td>
<td>10.0</td>
<td>0.0</td>
</tr>
<tr>
<td>$E_0$ (mV/m)</td>
<td>0.0</td>
<td>0.0</td>
<td>0.0</td>
<td>0.0</td>
<td>80</td>
<td>80</td>
<td>80</td>
<td>150</td>
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</tr>
</tbody>
</table>
Figure 3. The O⁺ temperature anisotropy factor and (b) field-aligned velocity from different combinations of DC electric fields and wave heating, at \( t_0 + 30 \) s, extracted from the center geomagnetic field line of the simulations. See Table 2 for a complete listing of the drivers used in each simulation. Increasing the wave power spectral density not only increases the anisotropy factor at high altitudes but also serves to increase the penetration depth of the wave heating effects into the ionosphere. If the wave heating co-occurs with a DC electric field, the DC electric field generates larger anisotropies at lower altitudes and can completely mask any low-altitude wave heating effects on temperature anisotropy.

velocities to a smaller degree at altitudes above the peak of the upflow (e.g. compare the different colors of each line type used in Figure 3b). When the wave heating is concurrent with a DC electric field, the electric field effects can completely mask any low-altitude wave heating effects on temperature anisotropy.

At altitudes above the main ion perturbation, a smaller level of upflow is generated by electrons that gain energy through heat exchange with frictionally heated ions deep in the ionosphere (i.e., near the F region peak). Electrons have a high thermal conductivity so any energy input at low altitudes is quickly conducted along the field line. Hence, an ion heat source can serve to generate, indirectly, electron pressure enhancements and ambipolar upflow in the topside ionosphere, here at 600–1000 km, before the main ion perturbation (seen in Figure 3b, at about 520 km altitude) can reach these altitudes. As an example of this effect, the simulation using only a DC electric field of 150 mV/m (case 9) has a field-aligned velocity \( \sim 60 \) m/s larger at 800 km, well above the main ion perturbation, than the simulation that does not use any upflow drivers (case I).

Figure 4 shows the time evolution of the O⁺ flux at 1000 km (Figure 4a) and 2500 km (Figure 4b) on the center geomagnetic field line of the 12 simulations summarized in Table 2. The rapid increase in ion flux, e.g., between

Figure 4. O⁺ flux versus time for different combinations of DC electric field and wave heating. (a) The O⁺ flux from the center geomagnetic field line of the model at 1000 km and (b) the flux at 2500 km (along the center geomagnetic field line of the model) for the 12 simulations listed in Table 2.
110 s and 180 s for case IX at 1000 km, in this figure indicates the arrival of the main ion perturbation primarily driven by frictional heating from DC electric fields. The wave heating only simulations, cases II, III, and IV, take a longer period of time for the ion flux response to build to full strength; a maximum response rate is not reached by the end of the simulation (10 min) at 2500 km. This may limit the impact on plasma supply to higher altitudes in absence of another upflow mechanisms, except in some extreme cases. When a DC electric field is included with wave heating effects, there is a definite increase in the flux response at 1000 km, and the flux maximum occurs more quickly. However, these strong fluxes decrease more rapidly than those generated by smaller DC electric fields or wave heating only situations.

Overall, the effects of wave heating plays a larger role at 2500 km than at 1000 km. At 2500 km, cases IV, VIII, and XII, which all used an extreme wave power spectral density of $|E_0|^2 = 10$ (mV/m)$^2$/Hz, consistently result in a larger flux than other ion driver combinations. An exception to this is the main ion perturbation arrival at 510 s of case XI, which is slightly larger than the ion flux response of the wave heating only simulation, case IV, at that point in time. At 1000 km the DC electric field plays a stronger role. Cases IX, X, XI, and XII all use an intense DC electric field of $E_0 = 150$ mV/m in addition to various levels of wave heating (cf. Table 2), and the resulting upward ion fluxes are the largest, and quickest, ranging in ion flux from $6.9 \times 10^{13}$ to $8.9 \times 10^{13}$ m$^{-2}$ s$^{-1}$. The indirect heat transfer from the ions to electrons has the effect of increasing ion flux at times before the arrival of the main ion perturbation. This indirect heat transfer mechanism can create a significant amount of ion flux, for example, $2 \times 10^{13}$ m$^{-2}$ s$^{-1}$ after 150 s at 1000 km for case V which does not contain any wave heating, only a DC electric field of 80 mV/m.

3.3. Thermospheric Wind Effectson Ion Upflow

Strong thermospheric (neutral) winds are fairly commonplace at high latitudes [e.g., Anderson et al., 2011], are known to play a role in regulating $F$ region ion dynamics, and have the potential to significantly impact the ionospheric upflow process. These winds would rarely be large enough to, alone, generate a large upflow. However, when coupled with other upflow mechanisms (e.g., DC electric fields and wave heating), winds may regulate upward ion fluxes and velocities. A sequence of simulations has been conducted using geographically horizontal neutral winds in different directions, in addition to DC electric fields and wave heating, to evaluate the degree to which winds may affect ionospheric upflow.

The first group of simulations, comprising eight model runs, illustrates the effects of wind in the geographic north-south direction. This set includes two reference simulations that lack any neutral winds, one that uses a moderate DC electric field only of $E_{0,\parallel} = 50$ mV/m (case II) and one that uses both a moderate DC electric field of $E_{0,\parallel} = 50$ mV/m and a typical wave heating with a power spectral density of $|E_0|^2 = 0.3$ (mV/m)$^2$/Hz (case VI). These two reference cases are compared to cases where horizontal neutral winds of $v_{n,0} = 100$ m/s southward, 200 m/s southward, and 100 m/s northward are individually added to these “base” upflow drivers. Table 3 lists the specific parameters used in each simulation. The neutral winds are specified as geographically horizontal and must be rotated into dipole coordinates (using equations (24) and (25)) prior to inclusion into the ion momentum and energy equations. In the Northern Hemisphere, southward winds have a component upward along the field line that induces upward plasma transport through drag and will tend to act synergistically with other ion upflow drivers included in the simulations. The northward winds have component downward along the field line and will tend to suppress ionospheric upflow.

Figure 5 plots the $O^+$ flux (Figure 5a) and field-aligned velocity (Figure 5b) along the center field line of the simulation domain after 10 min for the eight different combinations of applied drivers documented in Table 3. As shown in this figure, increasing the strength of southward neutral winds induces progressively larger $O^+$ field-aligned velocities and flux. A 200 m/s southward neutral wind exerts an upward force (through drag) that is enough to almost cancel the tendency for the $F$ region peak to drift downward due to pressure gradient.
and gravitational forces. Northward neutral winds exert a downward force enhancing the downflow under the \( F \) region peak and suppressing upflow at higher altitudes below the main ion perturbation, which has reached approximately 2000 km by the end of the simulation. In case V, downflow flux created by the northward neutral wind and upward flux driven by the wave heating balance resulting in the same amount of flux as the no wind, no wave heating simulation (case II) after 10 min at roughly 800 km. It is notable that through lifting of the \( F \) region plasma and through imparting parallel momentum to the ionosphere at lower altitudes dominated by ion-neutral collisions, neutral winds can significantly enhance upward flux and drift speeds at very high altitude regions.

A secondary effect of the neutral wind is through the frictional heating terms in equations (3) and (4). A second set of simulations was constructed to elucidate the impact of winds in the \( \mathbf{E} \times \mathbf{B} \) drift direction on ion upflow through the regulation of differential ion-neutral velocities, hence frictional heating. This set of simulations was conducted alongside a reference simulation that used both a DC electric field, \( E_{0 \perp} = 50 \text{ mV/m} \), and wave heating, \( |E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz} \) but lacked any neutral wind influence (case II, Figure 7 and Table 4). This reference case is compared against two cases, where horizontal neutral winds are used in conjunction with the reference simulation upflow drivers, \( E_{0 \perp} = 50 \text{ mV/m} \) and \( |E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz} \), 150 m/s along \( \mathbf{E} \times \mathbf{B} \) (case I, Figure 7 and Table 4) or 150 m/s against \( \mathbf{E} \times \mathbf{B} \) (case III, Figure 7 and Table 4). These winds are implemented using equation (26). Since there is not a component of these winds along the field line to cause ion-neutral drag, this wind orientation affects upflow through the frictional heating terms, in equations (2) – (4). As it can be seen in Figure 6, when the neutral winds are antiparallel to the \( \mathbf{E} \times \mathbf{B} \) drift (Figure 6c), then there is an increase in frictional heating which increases upflow, but when the neutral winds are parallel to the \( \mathbf{E} \times \mathbf{B} \) drift, then there is a decrease in frictional heating which suppresses upflow (Figure 6a).

**Table 4.** The Power Spectral Densities, \( |E_0|^2 \), DC Electric Fields, \( E_{0 \perp} \), and Geographic Neutral Winds, \( v_{n,0} \), Used in Simulations I–III Plotted in Figure 6*  

<table>
<thead>
<tr>
<th>Simulation</th>
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<th>III</th>
</tr>
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<tbody>
<tr>
<td>( v_{n,0} ) (m/s)</td>
<td>150 E</td>
<td>0.0</td>
<td>150 W</td>
</tr>
<tr>
<td>(</td>
<td>E_0</td>
<td>^2 ) (mV/m)²/Hz</td>
<td>0.3</td>
</tr>
<tr>
<td>( E_{0 \perp} ) (mV/m)</td>
<td>50</td>
<td>50</td>
<td>50</td>
</tr>
</tbody>
</table>

*Eastward winds here are parallel to the \( \mathbf{E} \times \mathbf{B} \) drift direction, and westward winds are antiparallel.

Figure 6 plots a snapshot of the \( O^+ \perp \) perpendicular temperature of these three simulations after 10 min of the applied drivers, the end of the simulation. The simulation using winds along \( \mathbf{E} \times \mathbf{B} \) is in Figure 6a, the reference simulation with no winds is in Figure 6b, and the simulation using winds against \( \mathbf{E} \times \mathbf{B} \) is in Figure 6c. Note that both the electric field boundary conditions and wind inputs have a Gaussian envelope in the direction perpendicular to the field lines which generates the central temperature structure seen in these panels of Figure 6. By the end of the
Figure 6. Perpendicular temperature versus altitude and meridional distance after 10 min of $E_0 \perp = 50 \text{ mV/m}$ and $|E_0|^2 = 0.3 \text{ (mV/m)}^2/\text{Hz}$ (a) with 150 m/s winds along $E \times B$, (b) with no winds, and (c) with 150 m/s winds antiparallel to $E \times B$. Temperature increases are localized in the center of the grid by constraining the electric potential boundary conditions using equations (23)–(26). At 400 km, the winds antiparallel to $E \times B$ have increased perpendicular temperatures by 11.5%, while the winds along $E \times B$ decrease them by 9.7% with respect to the control case.

Simulations, on the center geomagnetic field line, at 400 km, the against $E \times B$ winds (Figure 6c) increase the O\textsuperscript{+} perpendicular temperature by 11.5% while the along $E \times B$ winds (Figure 6a) decrease it by 9.7% from the no wind simulation due to the differences in frictional heating rates.

The resulting O\textsuperscript{+} flux from these three simulations is plotted in Figure 7. The along $E \times B$ wind simulation is case I, the no wind simulation is case II, and the against $E \times B$ wind simulation is case III. The against $E \times B$ neutral wind simulation not only increases the perpendicular temperature but also increases the O\textsuperscript{+} flux generated through the frictional heating mechanism. The leading edge of the primary upflow reaches 1000 km by 240 s (Figure 7a) and is still propagating toward 2500 km by the end of the simulation (Figure 7b). The flux at 2500 km is an order of magnitude smaller than the flux at 1000 km and is primarily driven by wave heating and the indirect heating mechanism whereby frictionally heated ions undergo heat exchange with electrons, which then transport the energy quickly along the field line.

A third and final group of simulations examine the dependence of the ion upflow response to neutral wind disturbance onset timing. It is highly unlikely that multiple ion drivers will occur at the exact same moment in time, so it is helpful to investigate how the relative timing of different energy source may affect ion upflow.

Figure 7. The O\textsuperscript{+} flux from three simulations at (a) 1000 km and (b) 2500 km along the center field line for the full 10 min duration of the simulation. For case descriptions see Table 4. Winds antiparallel to $E \times B$ (case III) increase the O\textsuperscript{+} flux response, while winds parallel to the $E \times B$ drift (case I) decrease the O\textsuperscript{+} flux response with respect to the reference simulation lacking winds (case II).
Figure 8. The $\text{O}^+$ (a) flux, (b) field-aligned velocity, and (c) density on the center geomagnetic field line, at 1000 km, for five simulations illustrating the effects of neutral wind onset time on the upflow responses. See Table 5 for a description of the different cases. Running neutral winds for a period of time prior to applying a DC electric field serves to increase densities and fluxes at higher altitudes but not significantly increase the field-aligned velocity.

Figure 8 shows the $\text{O}^+$ flux (Figure 8a), field-aligned velocity (Figure 8b), and the density (Figure 8c) for the last 10 min of five simulations used for this part of the study. For three of these simulations, first the model was run for an hour without any drivers being applied, then for case I $E_{\perp 0} = 50$ mV/m was active for 10 min (reference case), for case II $E_{\perp 0} = 50$ mV/m and a southward $v_{n,0} = 100$ m/s was active for 10 min, and for case III $E_{\perp 0} = 50$ mV/m and a southward $v_{n,0} = 200$ m/s was active for 10 min. For the next simulation the model was run for an hour with a constant southward $v_{n,0} = 100$ m/s and then $E_{\perp 0} = 50$ mV/m was also activated for 10 min in case IV. For the last simulation, the model was run for an hour with a constant southward $v_{n,0} = 200$ m/s and then $E_{\perp 0} = 50$ mV/m was also active for 10 min for case V. Hence, these five simulations vary in whether the neutral wind begins an hour before the main DC electric field or at the same time. For reference, the case parameters are also summarized in Table 5.

Running neutral winds for an hour prior to the DC electric field onset serves to increase ion densities at higher altitudes through ion-neutral drag. The 100 m/s southward neutral wind increases the $\text{O}^+$ density by 9% after one hour, and the 200 m/s southward neutral wind increases it by 20% at an altitude of 1000 km. These mark the initial state, $t = 0$, for Figure 8c. The resulting flux, by the end of the simulation, in case V is 22% larger, case IV is 11% larger, case III is 35% larger, and case II is 16% larger than the case without winds (case I). In either onset time scenario, given the parameters used here, larger neutral winds generate larger the $\text{O}^+$ flux responses as shown in Figure 8a. However, by starting the neutral winds an hour prior to the onset of a DC electric field, the resulting $\text{O}^+$ flux is lower than from the mutual onset case. Additionally, there is not an increase in the field-aligned velocity above that which is caused by the DC electric field alone. The sudden onset of a neutral wind disturbance has a larger impact on ion upflow, but in reality the neutral winds will act somewhere in between the two extremes simulated here, potentially taking tens of minutes to ramp up to speed if driven by magnetospheric energy inputs.

Table 5. The DC Electric Fields, $E_{\perp 0}$, and Geographic Neutral Winds, $v_{n,0}$, Used in Simulations I–V Plotted in Figure 8, Where $t_0$ Corresponds to 0 s

<table>
<thead>
<tr>
<th>Simulation</th>
<th>I</th>
<th>II</th>
<th>III</th>
<th>IV</th>
<th>V</th>
</tr>
</thead>
<tbody>
<tr>
<td>$v_{n,0}$ (m/s)</td>
<td>0</td>
<td>100°S</td>
<td>200°S</td>
<td>100°S</td>
<td>200°S</td>
</tr>
<tr>
<td>$v_{n,0}$ onset time (s)</td>
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<td>$t_0$</td>
<td>$t_0$</td>
<td>$t_0$-3600</td>
<td>$t_0$-3600</td>
</tr>
<tr>
<td>$E_{\perp 0}$ (mV/m)</td>
<td>50</td>
<td>50</td>
<td>50</td>
<td>50</td>
<td>50</td>
</tr>
<tr>
<td>$E_{\perp 0}$ onset time (s)</td>
<td>$t_0$</td>
<td>$t_0$</td>
<td>$t_0$</td>
<td>$t_0$</td>
<td>$t_0$</td>
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</tbody>
</table>

*In simulations IV and V, the neutral wind was started an hour prior to the onset of the DC electric field.
4. Conclusions and Future Work

A new anisotropic fluid model, GEMINI-TIA, based on a bi-Maxwellian distribution has been developed and used to study high-latitude ionospheric upflow driven by neutral winds, frictional heating, and wave heating processes. A comparison of this new model and its parent isotropic model shows that in situations of strong frictional heating, isotropic models may over estimate field-aligned velocities by up to 48%. This is a consequence of the use of a single pressure in the isotropic model which overestimates the impact of frictional heating on parallel transport. A comparison of plasma supplied to higher altitudes likewise shows significant differences in the model fluxes, indicating that it is likely important to resolve low-altitude temperature anisotropies for detailed modeling of ion upflow and outflow.

The new model has also been used to examine the synergistic effects of frictional heating and wave heating at low altitudes. At lower altitudes (<300 km) temperature anisotropies are largely driven by DC electric fields, while wave heating effects dominate above ~500 km. The strength of the power spectral density of broadband ELF waves determines how deep into the ionosphere these waves are able to generate anisotropy. Extreme levels of wave heating (e.g., power spectral density of 10.0 (mV/m)²/Hz) are required overcome the collisional relaxation and generate significant impacts at altitudes <300 km. This extreme power spectral density is much larger than the reference rate of 0.3 (mV/m)²/Hz taken from Bouhram et al. (2003) and would not be representative of a typical upflow/outflow event.

Neutral winds also play an important role in influencing ion dynamics. They can aid or hinder ion flow, and given enough time impart momentum to the ions at low altitudes which can impact the high-altitude ion populations available for secondary acceleration processes that lead to outflow to the magnetosphere. A geographically southward neutral wind of 100 m/s aids ion upflow, through ion-neutral drag, and can increase the O⁺ flux response by 15% at 1000 km after 10 min in a simulation that uses both that southward neutral wind and a DC electric field of 50 mV/m when compared to the DC electric field only simulation. On the other hand, geographically northward winds hinder ion upflow. A 100 m/s northward neutral wind coupled with a DC electric field of 50 mV/m results in an ion flux that is 13% smaller, at 1000 km after 10 min, than a similar simulation without winds. Neutral winds antiparallel to the E × B drift (westward here) exacerbate frictional heating, resulting in more ion upflow. A neutral wind against E × B of 150 m/s, coupled with a DC electric field of 50 mV/m and wave heating with a power spectral density of 0.3 (mV/m)²/Hz, increased the O⁺ flux response by 11.5% at 400 km, but a neutral wind along E × B of the same strength decreased the O⁺ flux response by 9.7% at 400 km after 10 min. The time history of neutral winds has also been shown to be important. By starting the neutral winds an hour prior to the onset of a DC electric field driven upflow, the resulting O⁺ flux is reduced by up to 11% in the case of 200 m/s southward winds by the end of the simulation when compared to cases when the neutral winds and DC electric fields are started at the same time. These results suggest that thermospheric dynamics can be an important factor affecting ion upflow and outflow.

Future work will include model validation activities that involve constraining the inputs with incoherent scatter radar (ISR) data and in situ rocket or satellite measurements. GEMINI-TIA is particularly well suited for interpreting multimodal measurements conducted as part of typical rocket campaigns (cf. Zettergren et al., 2014) or satellite campaigns with planned ground conjunctions. Data used to constrain the model can include perpendicular electric fields, characteristic energy and total energy flux from ISR, available neutral density and wind measurements, and precipitating electron fluxes. Results from these constrained simulations can then be compared against independent observations of the ion density and velocity responses. This allows us to evaluate ionospheric upflow response due to different auroral drivers having realistic temporal and spatial dependences.

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References


