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Electron and Ion Acceleration Associated with Magnetotail Reconnection

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Electron and Ion Acceleration

Associated with Magnetotail Reconnection

A dissertation submitted in partial satisfaction of the
requirements for the degree Doctor of Philosophy

in Physics

by

Haoming Liang

2017
This dissertation is dedicated to understanding electron and ion acceleration associated with magnetotail reconnection during substorms by using numerical simulations. Electron dynamics were investigated by using the UCLA global magnetohydrodynamic (MHD) model and large scale kinetic (LSK) simulations. The neutral line configurations and magnetotail flows modify the amounts of the adiabatic and non-adiabatic acceleration that electrons undergo. This causes marked differences in the temperature anisotropy for different substorms. In particular, one substorm event analyzed shows $T_\perp > T_\parallel$ ($T_\perp / T_\parallel \approx 2.3$) at ~$10R_E$ while another shows $T_\parallel > T_\perp$ ($T_\perp / T_\parallel \approx 0.8$), where $T_\perp$ and $T_\parallel$ (second order moments of the distribution functions) are defined with respect to the magnetic field. These differences determine the subsequent acceleration of the energetic electrons in the inner magnetosphere.
Whether the acceleration is mostly parallel or perpendicular is determined by the location of dayside reconnection.

A 2.5D implicit Particle-in-Cell simulation was used to study the effects produced by oxygen ions on magnetotail reconnection, and the associated acceleration of protons and oxygen ions. The inertia of oxygen ions reduces the reconnection rate and slows down the earthward propagation of dipolarization fronts (DFs). An ambipolar electric field in the oxygen diffusion region contributes to the smaller reconnection rate. This change in the reconnection rate affects the ion acceleration. In particular 67% of protons and 58% of oxygen ions were accelerated in the exhaust (between the X-point and the DF) in a simulation corresponding to a magnetic storm in which there was a 50% concentration of oxygen ions. In addition, 42% of lobe oxygen-ions are accelerated locally by the Hall electric field, far away from the X-point without entering the exhaust. Protons at the same locations experience $\mathbf{E} \times \mathbf{B}$ drift. This finding extends previous knowledge that oxygen and proton acceleration associated with reconnection mainly occurs in the exhaust and is consistent with Cluster observations. Oxygen ions and protons in the pre-existing current sheet are reflected by the DFs. The reflected oxygen beam forms a hook-shaped signature in phase space. In principle, this signature can be applied to deduce the DF speed history, and thus lead to remote-sensing of the reconnection dynamics.
The dissertation of Haoming Liang is approved

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PUBLICATIONS AND SELECTED PRESENTATIONS


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CHAPTER 1

Introduction

The aim of this dissertation is to investigate the electron and ion acceleration associated with reconnection in the Earth’s magnetotail. This chapter presents a background and literature review about magnetotail reconnection, geomagnetic activity, reconnection theory, plasma acceleration in the magnetotail and the involvement of oxygen ions in the geomagnetic activity. There are two factors that influence the magnetotail reconnection and thus the consequent acceleration of the electrons and ions: (1) the magnetotail configuration that results from the interaction of the magnetosphere with upstream solar wind conditions and (2) the presence of oxygen ions as a significant heavy ion species in the magnetotail. The motivation for each of the studies in this dissertation is discussed in Sections 1.4 and 1.5.

1.1. Reconnection in Earth’s Magnetosphere

In recent years there has been increased interest in magnetospheric physics during periods of extreme magnetic activity. This is motivated by the realization that severe magnetic activity can adversely affect communications, power networks and space exploration. That geomagnetic activity is driven by changes in solar wind conditions. The interplanetary magnetic field (IMF) and solar wind dynamic pressure are the two main factors that result in magnetospheric disturbances.

The interaction between the solar wind and the magnetosphere occurs at the dayside magnetopause. In 1961, Dungey [1961] introduced magnetic reconnection to explain a way in which the interaction between the solar wind and IMF interaction occurs. Magnetic
reconnection is a process by which oppositely directed magnetic field lines from the IMF merge with the Earth’s magnetic field. In this process, magnetic flux from the IMF is transferred to the magnetosphere leading to the transport of mass, momentum and energy from the solar wind to the magnetosphere. Figure 1-1 [Kivelson and Russell, 1995] presents a simple sketch of how the process works. It is a noon-midnight meridian cut of Earth’s magnetosphere. In this picture the solar wind brings southward IMF to the Earth. The IMF field line 1’ then reconnects with a closed geomagnetic line 1 at noon. After the reconnection, the resulting field lines are the open field lines labeled 2 and 2’ that are connected to both the Earth and sun. These opened flux tubes move continuously in time to the locations 3 and 3’, 4 and 4’, 5 and 5’ forming the lobe region in the magnetotail. As more and more open field lines accumulate in the lobes, the magnetic pressure compresses the plasma sheet and oppositely directed fields start another reconnection in the magnetotail as shown by 6 and 6’. In three dimensions, the reconnection sites are located along the cross-tail direction and form a line called an X-line or neutral line. The resulting field lines are labelled 7 and 7’. The field line 7’ returns back to the solar wind, while the field line 7 is a closed field line (closing in the Earth) and moves toward the Earth. It moves from location 7 to 8, bypasses the Earth, reaches location 9 returning flux back to the dayside. The motion of the points where the field lines intersect the ionosphere is indicated in the insert at the bottom of Figure 1-1. This flux tube motion is called convection.
1.2. Geomagnetic Activity

During relatively quiet times tail reconnection occurs in the distant tail (>60 \(R_E\)). However during moderate geomagnetic activity with prolonged (~30-60 minutes) southward IMF the reconnection occurs in the near-Earth tail and the steady reconnection becomes episodic. The input flux accumulates in the near-Earth tail and compresses the magnetotail plasma sheet leading to bursty reconnection in the region ~20-30 \(R_E\) away from Earth. This near-Earth X-line violently releases the stored magnetic energy and ejects fast-propagating hot plasmas and reconnected flux both earthward and tailward. The earthward propagating plasma
and flux further enhance the precipitation into the ionosphere and cause auroral disturbances.
As the lobe flux decreases, more and more closed field lines are generated and the increased pressure in the earthward plasma sheet pushes the near-Earth X-line tailward gradually. After all the excess magnetic energy in the lobes releases, the X-line will eventually reach the distant tail much as in quiet times. This sequence of events is called a magnetospheric substorm. The energy transfer in the magnetotail is fast with enhanced convection of magnetic flux and plasma to the inner magnetosphere.

In an isolated substorm, the perturbation of the magnetic field and the injection of plasma from the tail into the inner magnetosphere cause currents and precipitation in the ionosphere and lead to the activity in the aurora zone. If the strong southward IMF remains for a longer time (a few hours), there will be a prolonged period of energy input to the magnetosphere. Increased activity in the auroral zone will add ionospheric plasma including oxygen ions to the tail plasma sheet. In this case, the geomagnetic activity will be more complicated than in an isolated substorm and is referred as a geomagnetic storm. Energetic particles from the tail that are injected into the inner magnetosphere can undergo gradient and curvature drift motion in the dipole field configuration (the positive charged particles drift westward, while the negative charged particles drift eastward). These particles form a westward current called the ring current. The enhancement of the ring current generates large magnetic field disturbances opposite to the dipole field around the Earth and is a characteristic signature of a magnetic storm.

During storms, charged particles from the ionosphere fill the near-Earth plasma sheet. After they reach the tail plasma sheet, both the electrons and ions are accelerated by magnetotail reconnection and become hot species propagating earthward with reconnected magnetic flux and bulk flows. The earthward moving particles gain energy as they move in the electric and
magnetic fields. Although plasma is continuously injected into the inner magnetosphere, there are also various factors that can lead to its loss. Some of these particles are lost when they drift to the dayside and hit the eroded and compressed magnetopause. Some of them are scattered locally by the waves into the loss cone and precipitate into the ionosphere and some of the energetic ions charge-exchange with cold ionospheric ions, become neutral and are no longer confined by the magnetic field. When the solar wind input energy becomes weaker and reduces to the normal level, the charged particle injection rate will gradually become smaller than the loss rate. Then the ring current and related magnetic field disturbance will gradually decay. It takes a few days to recover to the normal level.

1.3. Theory of Magnetic Reconnection

Magnetic reconnection is a universal process occurring in many plasma phenomena, such as solar flares, laboratory plasmas, and magnetospheric substorms. It explosively converts large amounts of stored magnetic free energy into plasma bulk and thermal energy in a very short time, accompanied by changes of the magnetic topology. It is frequently demonstrated in empirically based models as a process of magnetic field line merging, breaking and reconnecting. [e.g., González and Parker, 2016]
In a plasma system, phenomena with large special scales (larger than the gyro-radii and inertial length of all species) and large temporal scales (larger than the cyclotron time, oscillation time and collision time of all species) can be well described by magnetohydrodynamic (MHD) equations. In MHD, the plasma is treated as an electrically conducting fluid. Figure 1-2 demonstrates the basic characteristics of fluid motion in MHD. Here, a fluid element is indicated by a loop drawn around it (the loop element is denoted by $d\vec{l}$). The black lines with arrows are the magnetic field lines. The magnetic flux passing through the fluid element can be indicated as $\int \vec{B} \cdot d\vec{S}$, where the $\vec{B}$ is the magnetic field and the $d\vec{S}$ is the area element in the direction of which can be found by applying the right-hand
rule to $d\vec{I}$. The change of magnetic flux in the frame of the moving fluid element can be expressed as

$$\frac{d}{dt} \int \vec{B} \cdot d\vec{S} = \frac{\partial}{\partial t} \int \vec{B} \cdot d\vec{S} + \oint \vec{B} \cdot \vec{V} \times d\vec{l}$$

(1-3-1)

Here, $\vec{V}$ is the velocity of the fluid element. On the right hand side, the first term is the local magnetic field change in time and the second term is the flux increase in the area swept by the loop. By using Faraday’s law and Stokes’ theorem, one can obtain

$$\frac{d}{dt} \int \vec{B} \cdot d\vec{S} = -\int \nabla \times \vec{E} \cdot d\vec{S} - \oint d\vec{l} \cdot \vec{V} \times \vec{B}$$

(1-3-2)

$$= -\oint d\vec{l} \cdot (\vec{E} + \vec{V} \times \vec{B})$$

For ideal MHD, Ohm’s law gives $\vec{E} + \vec{V} \times \vec{B} = 0$. This indicates that the magnetic flux is conserved on a moving fluid element and is said to be frozen into the flow. An ideal MHD fluid does not allow the breaking and reconnecting of the magnetic field lines; otherwise, the field lines will lose the identities of the specific fluid elements. Reconnection requires the diffusion of magnetic field lines in the fluid elements. This can be achieved when $\vec{E} + \vec{V} \times \vec{B} \neq 0$. One example is resistive MHD, in which the Ohm’s law is $\vec{E} + \vec{V} \times \vec{B} = \eta \vec{J}$, where $\eta$ is resistivity due to collision and $\vec{J}$ is current density. By substituting this Ohm’s law into Ampere’s law without high frequency oscillations and applying Faraday’s law, one can obtain

$$\frac{\partial \vec{B}}{\partial t} = \nabla \times (\vec{V} \times \vec{B}) + \frac{\eta}{\mu_0} \nabla^2 \vec{B}$$

(1-3-3)

Here, $\mu_0$ is the permeability of free space. On the right, the first term is the flux increase due to convection as shown in (1-3-2); the second term is the field line diffusion due to the resistivity.
The reconnection model for resistive MHD was proposed by Sweet and Parker [Sweet, 1958; Parker, 1957]. The Sweet-Parker Model is shown in Figure 1-3. This model assumes an antiparallel magnetic field configuration. It also assumes that the reconnected system reached a steady state, i.e., all the terms with \( \frac{\partial}{\partial t} \) vanish. The black lines are magnetic field lines and the gray box at the center is the region with field line diffusion in the fluid elements. The gray box with width \( 2\delta \) and length \( 2\Delta \) is called the diffusion region. Outside the gray box, the frozen-in condition is valid. During this steady-state reconnection, magnetic flux constantly flows toward the diffusion region along Z direction and after reconnection it is ejected away in X direction. The inflow velocity is \( V_{in} \) and the outflow velocity is \( V_{out} \). The mass density \( \rho \) is uniform. In the resistive MHD, \( \vec{E} = -\vec{V} \times \vec{B} + \eta \vec{J} \). The current density at the neutral sheet is in the Y direction since the magnetic field reversal is in the X direction, \( J \sim \frac{B}{\mu_0 \delta} \). In the inflow, the current density is small and the electric field is generated by the inflow motion of the
magnetic field $V_{in}B$. Near the reconnection point (at the center of diffusion region), i.e., the X-point, $\vec{B} \to 0$ and $\eta \vec{J}$ sustains the electric field. Because of the steady state, the electric field $\vec{E}$ is constant everywhere. Therefore, $V_{in}B \sim \frac{\eta B}{\mu_0 \delta}$, i.e., $V_{in} \sim \frac{\eta}{\mu_0 \delta}$. Due to the conservation of mass, $V_{in} \Delta = V_{out} \delta$. In MHD, the $V_{out}$ will eventually reach the Alfvén velocity $V_A$. Now, we have $\frac{\eta \Delta}{\mu_0 \delta} = V_A \delta$, i.e., $\frac{\delta}{\Delta} = \sqrt{\frac{\eta}{\mu_0 V_A \Delta}} = S^{-1/2}$, where $S$ is Lundquist number, which is defined as the ratio of the diffusion time to the Alfvén time. The most important quantity to evaluate the reconnection model is the reconnection rate. The reconnection rate is the magnetic flux generation rate in the outflow region (i.e., the integral of the generation rate of magnetic field in the outflow, $\frac{\partial B_z}{\partial t}$, over the neutral plane), which estimates how fast the reconnection releases the stored magnetic energy. By integrating the generation rate of the outflow magnetic flux and applying Faraday’s law and Stokes’ theorem, one can derive that the out-of-plane electric field $E_y$ at the X-point is equal to the reconnection rate. This electric field is called the reconnection electric field. Then we find the reconnection rate $E_y \propto V_{in} \propto S^{-1/2}$. However, in most of astrophysical systems, the Lundquist numbers are very large, which leads to very small reconnection rate compared to observations. As a result the Sweet-Parker model fails to explain the reconnection phenomena in such systems.
Figure 1-4 Configuration of collisionless reconnection. The thicknesses of ion diffusion region and electron diffusion region are in the scale of \( \frac{c}{\omega_{pi}} \) and \( \frac{c}{\omega_{pe}} \) respectively, where \( c \) is the speed of light, \( \omega_{pi} \) is ion plasma frequency, and \( \omega_{pe} \) is electron plasma frequency. (after Zweibel and Yamada [2009])

The plasma in the magnetotail is collisionless since the mean free path is much larger than the spatial scale of the magnetosphere. The resistivity should be very small (although some studies argue that scattering by the plasma waves and the turbulent electromagnetic field on the gyro-motion scale can produce anomalous resistivity [e.g., Labelle and Treumann, 1988; Cattell et al., 1995], this effect will not be discussed in this dissertation). Since the reconnection
occurs on diffusion scales in which the ideal MHD approximation is invalid, a mechanism beyond MHD needs to be invoked to explain the reconnection. A two-fluid model is a good option. In a two-fluid model, the electron and ion fluid equations are considered individually so that the spatial and temporal scale of their motion can be separated in the diffusion region. The thickness of a diffusion region with ion-scales (gyro-radius or inertial length) has been observed by spacecraft [e.g., Øieroset et al., 2001; Runov et al., 2006]. A schematic of reconnection site is shown in Figure 1-4. The magnetic field lines, current, ion flow, and electron flow are shown as black vectors, red vectors, blue dashed vectors, and black dashed vectors, respectively. The center is the X-point. The gray box is the ion diffusion region (IDR), where the ions are demagnetized, while the electrons are still frozen-in due to their much smaller gyro-radii. The white box inside the IDR is the electron diffusion region (EDR) where the electrons are demagnetized as well. The IDR and EDR scale with the gyro-radii and inertial lengths of ions and electrons respectively. Both the ions and electrons are frozen-in outside the IDR, where the MHD approximation is still valid. In the IDR, the ions are demagnetized and are diverted toward the outflow direction. In the EDR both species are unmagnetized and are diverted to the outflow. Note that the electrons near the separatrices are accelerated along the field lines toward the X-point before they reach the EDR [e.g., Yamada et al., 2010]. Because of the separation of the ion and electron flows, an in-plane Hall current is generated, which then generates the out-of-plane quadrupole Hall magnetic field. The Hall magnetic field is located in the outflow region near the separatrices (i.e., the boundary between inflow region and outflow region). This indicates the outflow field lines are warped in the out-of-plane direction. Because of the out-of-plane field line motion, the Hall electric field is generated near the separatrices [Drake et al., 2006] pointing toward the neutral sheet.
In the two-fluid model, considering $m_e \ll m_i$ and ignoring the resistive term, the generalized Ohm’s law can be written as

$$\frac{\hat{E} + \hat{V} \times \hat{B}}{n e} = \frac{1}{n e} \hat{J} \times \hat{B} - \frac{1}{n e} \nabla \cdot \hat{P}_e + \frac{m_e}{n e^2} \frac{\partial \hat{J}}{\partial t}$$  (1-3-4)

Here, $n$, $e$, $m_e$ are the number density, the unit charge, the electron mass and $\hat{P}_e$ is the electron pressure tensor. The non-zero terms on the right provide non-ideal conditions in MHD, which allow the diffusion of magnetic field lines. The terms on the right are the Hall term, the electron pressure tensor term, and the electron inertia term. In the IDR, the demagnetized ions do not move too much due to the large inertia, while the electrons keep moving with the field lines. Therefore, the Hall current is mainly carried by the electrons and the Hall term is significant in the IDR. In the EDR, $|\hat{B}|$ is very small as is the Hall term. There the other two terms become significant. Recent research has shown that the EDR can be divided into an inner EDR (characterized by an out-of-plane electron current due to the reconnection electric field) and the outer EDR (characterized by a super-Alfvénic electron outflow jet) [e.g., Shay et al., 2007; Karimabadi et al., 2007; Divin et al., 2012a]. If there are multiple X-points along X direction, the reconnected flux between two X-points will grow as a flux rope and form a magnetic island, or plasmoid. The spatial scale of magnetic islands can be larger or smaller than the diffusion region. As will be discussed later, magnetic islands are important for plasma acceleration.

1.4. Plasma Acceleration in the Magnetotail

The plasma acceleration in the magnetosphere during geomagnetic activity mainly occurs at the bow shock, in the magnetotail and in the radiation belts. In the magnetotail, the ions and
electrons are accelerated both in the vicinity of the X-line and during their earthward convection. In addition to the solar wind the plasma in the magnetotail mainly comes from the ionosphere outflow due to geomagnetic field perturbations, joule heating, and energetic particle precipitation. Along open field lines, ionospheric plasmas can reach the tail lobe with energies of a few tens of electronvolts (eV). They enter along closed field lines. The lobe ions and electrons can be transported into the plasma sheet through the plasma sheet boundary layer (PSBL) and reach energies of hundreds of eV to a few keV. They can also be involved and accelerated by magnetotail reconnection and then enter the plasma sheet as part of a reconnection outflow jet. Bursts of energetic electrons and ions with energies of a few hundreds of keV related to a magnetotail reconnection jet are observed by spacecraft [e.g., Baker and Stone, 1977; Möbius et al., 1983]. Recent studies observed the Hall field signature of magnetic reconnection accompanying these supra-thermal plasmas. The supra-thermal energy is much higher than the thermal energy and some studies suggest it is associated with the acceleration by reconnection. By using the Geotail satellite at \( X = -24R_E \), Hoshino et al. [2001] found that near a reconnection site, the supra-thermal electrons with energy \( E_e \geq 20 \) keV add a high energy power law tail to the thermal electron distribution. They proposed a two-step acceleration process: first, the acceleration occurs in the vicinity of the X-line; then further acceleration occurs downstream of the reconnection outflow. Øieroset et al. [2002] observed supra-thermal electrons with a few hundred of keV by using Wind spacecraft data in and around the ion diffusion region. Imada et al., [2007] observed \( E_e \sim 127 \) keV electrons near a magnetotail X-line which they explain by the two-step acceleration process. In addition, there are some observations suggesting that the acceleration of the electrons in magnetic islands is significant [e.g., Chen et al., 2008; Retinò et al., 2008; Wang et al., 2010; Huang et al., 2012a].
1.4.1. Plasma Acceleration near Magnetotail X-line

Both the ions and electrons are accelerated either adiabatically or non-adiabatically near the magnetotail X-line [e.g., Eastwood et al., 2013]. On one hand, when the scale of variation of the electromagnetic field is smaller than the scale of their gyro-motion, the particles tend to experience non-adiabatic acceleration e.g., typically in their diffusion regions. Such processes include the acceleration by the reconnection electric field near the X-point [e.g., Pritchett, 2006a, 2006b; Ricci et al., 2003; Lapenta et al., 2016], parallel electric fields at the separatrices generated by electron holes or the electron pressure tensor anisotropy [e.g., Drake et al., 2003, 2005; Egedal et al., 2005, 2012, 2013; Lapenta et al., 2014], Speiser motion or meandering motion in the region with $B_x$ reversal and significant $E_y$ [e.g., Speiser, 1965, 1967; Moses et al., 1993], and wave-particle interactions with a turbulent electromagnetic field, lower hybrid waves, kinetic Alfvén waves and whistler waves [e.g., Cairns and McMillan, 2005; Artemyev et al., 2009; Deng et al., 2009; Ashour-Abdalla et al., 2014; Fujimoto, 2014]. Ions can be accelerated when they are re-magnetized or picked up by the magnetic flux outflow jet [Drake et al., 2009; Lapenta et al., 2016] and they can form a hot distribution function with counter-streaming signature due to the trapping motion in the exhaust by the Hall electric field.

On the other hand, when the ions and electrons are transported into the region in which the scale of variation of the electromagnetic field is much larger than their gyro-scale, they will be able to experience adiabatic acceleration. In adiabatic acceleration the magnetic moment $\mu = \frac{W_\perp}{B}$, where $W_\perp$ is the perpendicular kinetic energy is conserved. The acceleration processes include the betatron acceleration due to travel from the weak magnetic field region to the magnetic field pile-up region [e.g., Fu et al., 2013; Huang et al., 2015], Fermi acceleration on the propagating outflow magnetic field structure, [e.g., Fu et al., 2013], Fermi
acceleration by contracting magnetic islands [Drake et al., 2006], gradient and curvature drift across $E_y$ in the outflow magnetic flux pile-up region [e.g., Hoshino et al., 2001]. Note that the betatron and Fermi acceleration near the reconnection site are associated with the reconnected magnetic flux pile-up region propagating away from the X-point. These mechanisms can occur simultaneously and keep accelerating the plasma when the reconnected flux is propagating earthward and much further away from the reconnection site. This will be discussed below.

### 1.4.2. Plasma Acceleration during Earthward Convection

During a geomagnetic substorm, closed plasma sheet magnetic field lines become stretched by the magnetotail reconnection process. Due to the strong magnetic tension, the reconnected lines tend to shrink toward the Earth and become more and more dipole-like. This process is called dipolarization. The magnetic structure of the dipolarization is also referred as a bubble [Wolf et al., 2009] and is characterized by the enhanced magnetic field, high temperature, low density, low plasma pressure and low specific entropy. Typically, the field lines of dipolarization pile up at the earthward edge due to the pressure balance with the preexisting plasma sheet in front of it. This pile-up region of the dipolarization is called a dipolarization front (DF). Note that an observational definition of a DF is a sudden enhancement of $B_z$ observed by spacecraft frequently preceded by a decrease in front of the enhancement. It only indicates the pile-up of $B_z$ due to any compression in the tail but does not mean that the DF definitely comes from a reconnection site. DFs have been explained as reconnection jet fronts [e.g., Sitnov et al., 2009; Runov et al., 2012], interchange-generated flow heads [e.g., Pritchett and Coroniti, 2011], or a consequence of the ion tearing mode [e.g., Sitnov et al., 2013]. Observations show that a DF is about an ion inertial length or ion gyro-radius thick when
it is located between $X= -11 \, R_E$ and $-20 \, R_E$ [e.g., Runov et al., 2009]. The Hall current is also found in the DF and scales similarly [e.g., Zhang et al., 2011; Fu et al., 2012]. Commonly, DFs are observed to accompany bursty bulk flows (BBFs), which are high-speed (hundreds of km/s) earthward flows which last one minute or so. Such high-speed flows are believed to be the major form of magnetic flux and plasma transport from the tail to the Earth [Angelopoulos et al., 1992, 1994]. Recent studies also imply that the BBFs are associated with the near-Earth X-line around $X= -20 \, R_E$ during a substorm [e.g., Angelopoulos et al., 2008].

Enhanced high energy fluxes of electrons and ions are frequently observed by spacecraft during DF events embedded in the earthward propagating BBFs [e.g., Runov et al., 2009; Zhou et al., 2009; Deng et al., 2010; Fu et al., 2011]. The electron fluxes from a few keV to a few hundred keV are enhanced (a few tens of keV to a few hundred of keV for ions) when a DF passes a spacecraft, while the fluxes below these energies decrease when the DF passes. The pitch angle of the enhanced flux is around 90 degrees in some events while it is close to 0 and 180 degrees in other events [e.g., Zhou et al., 2009; Hwang et al., 2011; Fu et al., 2011]. There are many acceleration mechanisms proposed to explain the energetic particles associated with the earthward propagating DFs and BBFs. They include betatron acceleration due to the transport of plasma from a region with weak magnetic field to a region with strong magnetic field [e.g., Zelenyi et al., 1990; Ashour-Abdalla et al., 2011; Fu et al., 2011; Pan et al., 2012], Fermi acceleration due to particles bouncing on earthward contracting magnetic field lines [e.g., Zelenyi et al., 1990; Fu et al., 2011; Pan et al., 2012; Ashour-Abdalla et al., 2013], acceleration by the cross-tail electric field during trapping between reversed magnetic fields with Speiser motion [e.g., Speiser, 1965, 1967], wave-particle scattering by the waves observationally associated with DFs, e.g., lower hybrid waves, electron cyclotron harmonics and whistler waves.
[e.g., Deng et al., 2010; Zhou et al., 2009; Huang et al., 2012b], and turbulent acceleration in the magnetotail plasma sheet [e.g., Zelenyi and Milovanov, 2004]. Fu et al. [2011] suggest that the different pitch angle dependences in the DF events are due to either betatron or Fermi acceleration.

The high energy plasma propagating with the DFs and BBFs is eventually injected into the inner magnetosphere in the flow braking region between geosynchronous orbit, ($X=-6.6 \, R_E$), and around $X=-10 \, R_E$ [e.g., Sergeev et al., 2009]. The injection distribution in this region shows dispersionless injection signatures, i.e., the fluxes in all the energy channels increase simultaneously. This fact indicates that the acceleration during the propagation of the plasma in the magnetotail is rarely energy-dependent and that the contribution of gradient and curvature drift along the cross-tail electric field is weak in the magnetotail compared to the energy gain from the betatron and Fermi acceleration. In the injection region, the field lines are more dipole-like and the ions and electrons start to have significant gradient and curvature drift along the cross-tail electric field. After they drift around the Earth back to the injection region several times, their distribution function shows multiple injection signatures with gradual dispersion, which is referred to as a drift echo [e.g., Lanzerotti et al., 1967].

1.4.3. Motivation of Studying Electron Acceleration in Realistic Tail Configurations

During substorm injection, the ions and electrons are accelerated in the magnetotail near the X-line and then in their convection earthward they become seed populations for the subsequent acceleration in the inner magnetosphere [e.g., Baker et al., 1979]. Many factors can influence the particle acceleration process near the magnetotail X-line and during the earthward convection. They include the configuration of the magnetotail and the highly dynamic X-lines, the variation of the parameters of DFs and BBFs, the generation of plasmoids [e.g., Hones,
1979], and plasma sheet flapping [e.g., Lui et al., 1978; Sergeev et al., 2003; Runov et al., 2005; Sitnov et al., 2014] and buoyancy [e.g., Roux et al., 1991; Panov et al., 2012a, 2012b; Sitnov et al., 2014].

However, many studies of acceleration mechanisms in the magnetotail are based on models with ideal magnetic field configurations [e.g., Zelenyi et al., 1990; Ashour-Abdalla et al., 1990, 1994; Schriver et al., 1998; Li et al., 1998; Birn et al., 2011, 2013]. The transport and acceleration mechanisms based on magnetotail configurations during different substorms are not fully understood. Since substorms can vary significantly from event to event due to the upstream solar wind conditions, it is important to perform simulations to study acceleration quantitatively based on realistic satellite events and realistic magnetotail configuration. As an example, the February 15, 2008, THEMIS event [Zhou et al., 2008] shows perpendicular energetic electrons behind the DFs, while the August 15, 2001, Cluster event [Hwang et al., 2011] shows the field-aligned electron distributions behind the DFs. The difference may be due to different contributions of betatron and Fermi acceleration, but it could also result from the configuration of the magnetotail and the X-lines.

Therefore, one of the outstanding questions is, how does the configuration of realistic magnetotail and X-lines affect the electron velocity distribution functions near the reconnection site and during their earthward convection? This question can be addressed by using UCLA global MHD [e.g., Raeder et al., 1995, 1998, 2001; El-Alaoui, 2001; El-Alaoui et al., 2009] accompanied by the test particle tracing technique called Large Scale Kinetics (LSK) [e.g., Ashour-Abdalla et al., 1990; Schriver et al., 1998]. The realistic magnetotail configuration can be reproduced by the global MHD model driven by the realistic solar wind conditions [e.g., Ashour-Abdalla et al., 2009, 2011; Pan et al., 2014a, 2014b]. To validate the simulation results,
the MHD+LSK simulation scheme is used to compare with the satellite observations of the magnetic field, plasma flows, particle distribution functions, and particle energy flux in different channels. This will be discussed in detail in Chapter 3 of this dissertation.

1.5. Oxygen Ion Circulation in the Magnetosphere

The heavy ions, mostly oxygen ions (O\(^+\)), are frequently observed throughout the magnetosphere [e.g., Shelley et al., 1972; Sharp et al., 1981; Mukai et al., 1994; Hamilton et al., 1988] during geomagnetic activity, especially geomagnetic storm. Understanding the physics of O\(^+\) in the magnetosphere is indispensable for space weather modeling. On one hand, although the heavy ions are found mainly originating from the ionosphere as a cold plasma source (a few eV), they are energized to keV even MeV in the plasma sheet and the inner magnetosphere accompanying geomagnetic disturbances. Because they have relatively large gyro-radii and gyro-period, they perform non-adiabatic motion most of time in the magnetosphere. On the other hand, the heavy ions change the typical plasma parameters like the mass and Alfvén speed and provide addition spatial and temporal scales, which can result in the variation of various physical processes in the magnetosphere, e.g., magnetic reconnection, tearing instability, Kelvin-Helmholtz instability, ion-scale wave generation and resonances, etc. The non-adiabatic motion of heavy ions also introduce more complicated non-linear evolution of these processes, e.g., current sheet bifurcation [e.g., Zelenyi et al., 2002; Cai et al., 2008]. They can thus globally change the threshold for substorm onset, the speed of magnetospheric convection, the level and duration of the geomagnetic disturbances, and the plasma energization during these processes.
Figure 1-5 Oxygen ion circulation in the magnetosphere (Adapted from [Kronberg et al., 2014])

Figure 1-5 shows the circulation of heavy ions in the magnetosphere. Due to the regular low abundance of heavy ions (~3%-6% for He\(^{++}\), ~0.1% for other heavy ions) in the solar wind [e.g., Neugebauer, 1981; Aellig et al., 2001], the heavy ions in the magnetosphere mainly come from the ionosphere outflow. Although various heavy ions are observed escaping from ionosphere, e.g., He\(^+\), N\(^+\), NO\(^+\) and O\(_2\)^+\, O\(^+\) is dominant and thus the physical process and consequences of only ionospheric O\(^+\) are considered in this dissertation.
1.5.1. Ionospheric O⁺ Outflow

The O⁺ outflow source regions are mainly the polar cap region (corresponding to open field lines) and the auroral region (corresponding to closed field lines at high latitudes). In the polar cap, the ions with low energy (a few eV) can escape due to the pressure gradient along the open field lines from the ionosphere to the lobes in the magnetotail and form a steady ion outflow, i.e., the polar wind. Because of their much smaller mass, the electrons have a larger scale height in the ionosphere than the ions and the separation leads to an ambipolar electric field. This ambipolar electric field can further accelerate the ions to higher altitudes. In the auroral region, the O⁺ ions are accelerated in the dayside cusp region (a few eV to a few keV) and in the nightside auroral region (up to a few tens of keV) during geomagnetic activity. The energy inputs are mainly the downward propagating Poynting flux from the reconnection at the dayside magnetopause and in the magnetotail plasma sheet and electron precipitation. They heat the ambient plasma at low altitudes, increase the ionization and electron scale height, and thus enhance the ambipolar electric field. At altitudes higher than 1000 km, the O⁺ ions are further accelerated by resonances with various waves [e.g., Andre and Yau, 1997 and references therein]. Due to the transverse heating by double layers [Borovsky, 1984], lower hybrid waves [e.g., Palmadesso et al., 1974], ion cyclotron waves [e.g., Ashour-Abdalla and Okuda, 1984], ion-ion hybrid waves [e.g., Schrifer and Ashour-Abdalla, 1988], and waves in the plasma sheet boundary layer (PSBL, i.e., open/closed field line boundary) [e.g., Tung et al., 2001], the ions obtain enough perpendicular energy and start to be accelerated along the magnetic field lines by the mirror force [e.g., Strangeway et al., 2005], which leads to commonly observed ion conics [e.g., Sharp et al., 1977; Shelley et al., 1976; Yau et al., 1983; Klumpar, 1979; Gorney et al., 1981]. By following the magnetic field lines, the O⁺ ions from the cusp can convect
tailward to the polar cap region and escape along the open field lines into the lobes in the magnetotail, while the O$^+$ ions from the nightside auroral region follow the closed field lines to the plasma sheet in the magnetotail and the inner magnetosphere.

1.5.2. Ionospheric O$^+$ Ions in the Lobe and the Plasma Sheet

The O$^+$ ions propagating in the lobes have parallel velocity along the field lines and perpendicular $\vec{E} \times \vec{B}$ convection toward the plasma sheet. The parallel speed determines how far away the tailward locations are when they enter the plasma sheet with the field lines. Therefore, the ions with higher speed can reach further locations from the lobe to the plasma sheet [e.g., Seki et al., 1998a, 1998b]. This is called the velocity filter effect. This effect explains the observations of the mono-energetic distribution of the O$^+$ ions in the polar cap [e.g., Elliott et al., 2001] and in the lobes [e.g., Lennartsson, 1994; Seki et al., 1998a]. By examining the distribution function in the cusp and assuming Liouville’s theorem, observations [e.g., Seki et al., 1998b; Liao, 2011] suggest that besides the velocity filter effect, additional acceleration is needed to explain the O$^+$ energy distribution in the tail lobes during storm times. Centrifugal acceleration has been suggested as the main acceleration mechanism for the ions in the lobe [e.g., Nilsson et al., 2010]. In addition, spacecraft observations show that before convection to the polar cap region, the O$^+$ ions at the high altitudes in the cusp can be heated significantly (up to ~10 keV) in the transverse direction by wave-particle interactions [e.g., Bouhram et al., 2004; Nilsson et al., 2006] and obtain much more energy than they at the low altitudes. The combination of these acceleration mechanisms is necessary to explain the energy of the O$^+$ ions when they enter the plasma sheet from the lobes in the near-Earth tail (~20 RE) and distant tail (~200 RE).
The ionospheric O+ ions reach the magnetotail plasma sheet by either following the closed field lines from the nightside auroral region or convecting through the PSBL from the lobe or pass through the magnetotail reconnection regions from the lobe. It is worth noting that the O+ density and O+/H+ density ratio in the plasma sheet are proportional to solar extreme ultraviolet radiation (EUV) [e.g., Young et al., 1982] and geomagnetic activity, as measured by the Auroral Electrojet (AE) index [e.g., Lennartsson and Shelley, 1986], the disturbance storm time (Dst) index [e.g., Bouhram et al., 2005], and Kp-index [e.g., Mouikis et al., 2010]. The solar EUV is mainly related to the solar cycle and solar activity. It increases the ionization and heats the O+ ions in the ionosphere. The AE index represents the geomagnetic perturbations during substorms with one-minute cadence due to the field-aligned currents connecting the tail plasma sheet and the ionosphere. The Dst index shows hourly geomagnetic disturbances resulting from the ring current enhancement, which is related to the storm activity. Kp measures the magnetic perturbation at mid-latitudes with relatively low resolution (3 hours). In addition, observations show that in the tail plasma sheet near the Earth (e.g., 7-8 RE), the solar EUV and geomagnetic activity have roughly equal contributions to the increase of the O+ density and the O+/H+ density ratio, while in the plasma sheet further away (e.g., 15-20 RE) geomagnetic activity is the dominant factor [Maggiolo and Kistler, 2014]. Both the O+ density and O+/H+ density ratio increase toward the Earth.

1.5.3. Acceleration of O+ Ions in the Plasma Sheet

In the plasma sheet, the O+ ions with energy from tens of keV to hundreds of keV have been observed. Near the PSBL, the O+ ions are heated to around a few keV due to the fluctuations caused by the Kelvin-Helmholtz instability [e.g., Grigorenko et al., 2010]. Unlike the protons and electrons, the O+ ions mainly experience non-adiabatic acceleration in the
plasma sheet, due to their much larger gyro-radii and much longer gyro-period. Because of the reversal of the stretched field lines in the tail plasma sheet, the O\(^+\) ions are energized when they perform Speiser motion [Speiser, 1965, 1967] or meandering motion across the convection electric potential [e.g., Lyons and Speiser, 1982; Ashour-Abdalla et al., 1994]. The O\(^+\) ions can also be accelerated by the inductive electric field accompanying the various transient structures or processes in the magnetotail, including the magnetic reconnection [e.g., Wygant et al., 2005], DF [e.g., Sharma et al., 2008], plasmoid release [e.g., Wilken et al., 1995; Zong et al., 2004], and turbulent electromagnetic fields [e.g., Milovanov and Zelenyi, 2002]. These can lead to mass-dependent ion acceleration. When the scales of the transient structures and processes are located between the spatial and temporal scales of the O\(^+\) gyro-motion and those of the H\(^+\) gyro-motion, the O\(^+\) ions are preferentially accelerated non-adiabatically by the impulsively inductive electric field, while the H\(^+\) ions may only perform the adiabatic convection [e.g., Möbius et al., 1987; Nosé et al., 2000]. Once the O\(^+\) ions inject into the ring current and the inner magnetosphere, they are energized not only by adiabatic gradient and curvature drift along the convection electric field but also by the non-adiabatic interaction with the low frequency electromagnetic fluctuations related to injected dipolarization and ion-scale waves [e.g., Nosé et al., 2010; Nosé et al., 2014].

1.5.4. Effects of the Ionospheric O\(^+\) Ions on the Magnetospheric Dynamics

O\(^+\) ions have significant effects on magnetospheric dynamics during enhanced geomagnetic activity, which result in pronounced differences between the proton-electron cases on both global and local scales. Globally, the ionospheric O\(^+\) ions increase the mass density of the plasma sheet and build up pressure in the tail which leads to considerable stretching of the magnetotail [e.g., Glocer et al., 2009a, 2009b]. Both the observations [e.g., Winglee, 2002] and
simulation [e.g., Glocer et al., 2009a] show that the increase of O\(^+\) outflow is accompanied by a reduction of the cross polar cap potential, which slows down the convection in the plasma sheet and the ring current, while the impact of the O\(^+\) outflow on ring current development is still debated [e.g., Fok et al., 2011; Welling et al., 2011]. Locally, when unmagnetized in the region with reversed field lines in the plasma sheet, the particles are on meandering orbits. Because of their large gyro-radii, the current sheet associated with the O\(^+\) ions is thicker than that associated with the protons and electrons [e.g., Zelenyi et al., 2006]. Although this has been confirmed by some observations [e.g., Cai et al., 2008; Petrukovich et al., 2011], Liu et al. [2014] found that the current sheet thickness is related to the proton gyro-radii instead of the gyro-scale of the O\(^+\) ions. In addition, the O\(^+\) meandering motion between the reversed field lines can result in the current sheet bifurcation [e.g., Ashour-Abdalla et al., 1994; Zelenyi et al., 2002; Cai et al., 2008].

The O\(^+\) ions can increase the linear tearing growth rate and reduce the threshold of tearing in the magnetotail [Baker et al., 1982]. This theory has been used to explain the dawn-dusk asymmetry of the location for the substorm onset [e.g., Frey et al., 2004]. However, the O\(^+\) density observations in the tail plasma sheet do not show such asymmetry [e.g., Mouikis et al., 2010; Maggiolo and Kistler, 2014]. By contrast, by observing the total pressure in the lobes and the plasma sheet, Liu et al. [2013] show that the maximum accumulated magnetic flux in the lobes before substorm onset is proportional to the O\(^+\) density and O\(^+\)/H\(^+\) density ratio, which indicates that the substorm is harder to trigger when there are more O\(^+\) ions. It is worth noting that the O\(^+\) ions are also suggested as an important factor that causes the sawtooth events. A sawtooth event is named for the periodic injection of high energy ion flux around geosynchronous orbit, which results from a sequence of substorms [e.g., Brambles et al., 2011;
This is explained as a consequence of a magnetosphere-ionosphere coupling loop in which the O\(^+\) outflow is enhanced due to the energy input of the substorm perturbation, and then builds up the pressure in the inner magnetosphere and stretches the tail field lines, which further causes the formation of new near-Earth neutral lines [e.g., Wiltberger et al., 2010; Ouellette et al., 2013].

The O\(^+\) ions also affect the electromagnetic ion cyclotron (EMIC) waves frequently observed in the ring current. EMIC waves are excited due to ion anisotropy [e.g., Cornwall, 1965] and by scattering and resonantly interacting with electrons and ions, they contribute to the loss of energetic particles in the radiation belts. They can also heat the ions due to wave-particle resonances at specific frequencies. With multiple ion populations, the dispersion relation of the EMIC waves is modified and consists of multiple bands related to the multiple ion gyro-frequencies [e.g., Hu et al., 2010; Chen et al., 2011]. Simulations [e.g., Hu et al., 2010; Omidi et al., 2013] show that when the fraction of O\(^+\) ions is less than 7\%, the EMIC waves can propagate along the field lines from the equator to be absorbed near the ionosphere, while with a moderate fraction of O\(^+\) ions (e.g., ~15\% in Omidi et al. [2013]) the EMIC waves are reflected at high latitudes at the bi-ion frequency (ion-ion hybrid frequency) and with a high fraction of O\(^+\) ions (e.g., >30\% in Omidi et al. [2013]) they can lead to strong ion heating.

1.5.5. Motivations of Studying the O\(^+\) in the Magnetotail Reconnection

Based on the above discussion, the ionospheric O\(^+\) ions have profound effects on the magnetospheric dynamics especially during enhanced geomagnetic activity. One of the most important magnetospheric processes is magnetotail reconnection. It not only heats the plasma but also controls the amount of released and the release rate of the stored magnetic energy in the tail. Therefore, in order to better model space weather, the top priority is to understand how
the ionospheric O\(^+\) ions affect the magnetotail reconnection and how the O\(^+\) ions are energized by the reconnection.

In spacecraft observations near the magnetotail X-line close to \(X = -19\) \(R_E\), the O\(^+\) concentration varies between storm-time and non-storm substorms [e.g., Kistler et al., 2005, 2006]. There is only a small fraction of O\(^+\) ions observed during non-storm substorms, while in some storm-time substorm events the O\(^+\) ions are observed to dominate the ion number density and pressure. The velocity distribution functions of the O\(^+\) ions show evidence of non-adiabatic meandering motion in the thin current sheet, which contributes 5%-10% of the cross-tail current. Luo et al. [2014] observed energetic O\(^+\) ions with energy of >150 keV associated with the formation of the magnetotail X-line. Some observations show that there are energetic/heated O\(^+\) ions in both earthward and tailward jets during reconnection [Zong et al., (1998)], in the plasmoids [e.g., Wilken et al., 1995; Zong et al., 1997], and in the vicinity of the X-point [e.g., Wygant et al., 2005].

There are a few simulation studies that try to investigate the effects of O\(^+\) on magnetotail reconnection and the O\(^+\) acceleration during the reconnection. Based on 2D kinetic simulations with periodic boundary conditions, Hesse and Birn [2004] found that the O\(^+\) ions do not have much effect on the reconnect rate, though cautioned readers that the limited size of the simulation box and limited duration of the simulation limited O\(^+\) motion. Shay and Swisdak [2004] analyzed three fluid equations and concluded that the presence of the O\(^+\) ions can result in a “heavy whistler” mode to mediate the reconnection and a “heavy Alfvén” mode to determine the reconnection rate. As demonstrated in their 2D multi-fluid simulation with periodic boundary conditions, in the presence of the O\(^+\) ions, the reconnection rate became smaller and the scale of the Hall magnetic field became broader than the case without O\(^+\) ions.
Using an implicit Particle-in-Cell (PIC) simulation with periodic boundary conditions, Markidis et al. [2011] studied 2D reconnection including $e^-$, $H^+$ and $O^+$ with realistic mass ratios. By initially setting a proton-electron Harris current sheet [Harris, 1962] with lobe $O^+$ ions, they show that the presence of the $O^+$ ions slightly reduces the reconnection rate and that the Hall magnetic field has a two-scale structure in the direction normal to the current sheet: a narrow sharp peak of the out-of-plane magnetic field related to the $H^+$ diffusion region and a broad slow decrease of the out-of-plane magnetic field related to the $O^+$ diffusion region. They also found that the $H^+$ and $O^+$ separate in the outflow region due to different acceleration mechanisms resulting from different spatial and temporal gyro-scales of the two ion species. Moreover, they found the presence of the $O^+$ ions does not change the characteristic signatures of the reconnection with a guide field. Karimabadi et al. [2011] used a linear Vlasov solver and 2D PIC simulations with driven and non-driven boundary conditions to study the $O^+$ effects on the tearing instability and the reconnection. They found that the linear growth rate of the tearing mode is only sensitive to the $O^+$ current carrier rather than the lobe $O^+$ ions. With an initial proton-electron Harris current sheet, the lobe $O^+$ ions replace the initial $H^+$ current and become a current carrier only after they enter the current sheet during the non-linear stage of the reconnection. As a result, the lobe $O^+$ ions reduce the energy conversion rate and change the composition of the current sheet. They also found the presence of the $O^+$ ions reduces the number and repetition frequency of the secondary islands and slows down the coalescence process. Using 2.5D PIC simulation with open boundary conditions, Liu et al. [2015] studied the acceleration of the $O^+$ ions from the lobe during the magnetotail reconnection. By calculating the agyrotropy [Scudder and Daughton, 2008], they identified the $H^+$ and $O^+$ diffusion regions and investigated the ion acceleration inside these regions. The velocity
distribution functions of the O$^+$ ions in the exhaust show a beam in the out-of-plane direction and counter-streaming signature in the direction normal to the current sheet. They found the similar signatures in two reconnection events observed by Cluster spacecraft. Using test particle simulations, they found that in the vicinity of the X-point, both H$^+$ and O$^+$ ions are accelerated by the electric field pointing in the outflow direction, while O$^+$ ions also obtain equal amount of energy from the out-of-plane reconnection electric field.

Although there were many studies focusing on the interaction between the O$^+$ ions and magnetotail reconnection, there are still some outstanding questions:

(1) Previous simulation studies mainly tried to address the issues related to the lobe O$^+$ ions only. However, some observations [e.g., Kistler et al., 2005] and simulations [e.g., Karimibadi et al., 2011] show that the O$^+$ ions can become a current carrier near the reconnection site. The question is, what effect do they have on the evolution of reconnection and reconnection-related structures, e.g., the reconnected jet fronts or the DFs?

(2) In MHD, outside the diffusion region, the O$^+$ ions reduce the Alfvén speed and thus reduce the reconnection rate. The reconnection rate measures how fast the magnetic flux transfer through the X-point is. However, there are no studies examining whether the O$^+$ ions in the O$^+$ diffusion region can slow down or accelerate or have an effect on magnetic flux transfer around the reconnection X-point. Therefore, the question is, what is the O$^+$ effect on the magnetic flux transfer in the O$^+$ diffusion region? This is important because: (a) the knowledge of such heavy ion effects can improve our understanding of kinetic reconnection beyond the MHD regime; (b) particle acceleration of protons and heavy ions near the reconnection site is closely related to
the inductive electric field associated with the magnetic flux transfer process. In a plasma with O\(^+\) the influence on particle acceleration may be significant, because based on Liu et al. [2015], the O\(^+\) diffusion region is much larger region than those for the other two species.

(3) According to the previous studies, the presence of the O\(^+\) ions can cause a broader Hall magnetic field structure [e.g., Shay and Swisdak, 2004; Markidis et al., 2011; Karimabadi et al., 2011]. How do the O\(^+\) ions affect the spatial scales of the reconnection-related regions/structures besides the Hall magnetic field, e.g., the outflow and the DFs? How does any effect change with the O\(^+\) concentration?

(4) Liu et al. [2014] show that the current sheet thicknesses near the observed reconnection sites are independent of the O\(^+\) scales but mainly related to the proton gyro-radii. Since the O\(^+\) ions have large gyro-radius compared to the thin current sheet thickness, they are likely involved in the reconnection mainly in a non-adiabatic way. Unlike the H\(^+\) ions, the O\(^+\) ions can non-adiabatically enter all the reconnection-related regions, e.g., the pre-existing current sheet, the exhaust, the DF region, the region with strong reconnection electric field near the X-point, the region with Hall fields near the separatrices, etc. The many different entry locations of the O\(^+\) ions may strongly influence their acceleration history. Thus the questions are, how do the O\(^+\) entry locations affect their acceleration processes and how different are the accelerated plasma states between the O\(^+\) and H\(^+\) in general after experiencing reconnection?

(5) Proton acceleration near the reconnection site includes both adiabatic and non-adiabatic processes [e.g., Drake et al., 2009; Lapenta et al., 2016]. How is proton acceleration changed in a plasma with O\(^+\) by all the effects mentioned above?
Because the evolution of the reconnection and its related structures are sensitive to the plasma states including electrons, protons, and oxygen ions, and vice versa, all the outstanding issues above will be addressed in this dissertation in a self-consistent way with a 2.5D implicit PIC simulation and self-consistent particle tracing.

1.6. Structure of the Dissertation

Chapter 1 has introduced the theory and observations of the magnetotail reconnection, the ionospheric plasma circulation, and the acceleration in the magnetosphere during geomagnetic activity. Chapter 2 introduces the numerical simulation methods that are used in the studies in this dissertation, including large scale kinetics (LSK) and implicit PIC simulation. Chapter 3 studies the non-adiabatic electron acceleration near the magnetotail X-line associated with different magnetotail configurations during substorms by using the MHD+LSK simulation scheme and observations. Based on implicit PIC simulations, Chapters 4 and 5 investigate the O$^+$ effects on the magnetotail reconnection and ion acceleration by reconnection in the presence of the O$^+$ ions. Chapter 6 presents a preliminary observation event that is related to the lobe O$^+$ acceleration associated with the magnetotail reconnection which is comparable to what was found in Chapter 5. Chapter 7 summarizes the findings and discusses unsolved problems and future research directions.
CHAPTER 2

Simulation Methods

This chapter introduces the two simulation methods used in the studies of this dissertation. First, the large scale kinetic (LSK) simulation is used to study electron acceleration in realistic magnetotail configurations during substorms (Chapter 3). Second, the implicit Particle-in-Cell (PIC) simulation is used in the studies of the effects of oxygen ions on magnetotail reconnection (Chapter 4) and acceleration of the oxygen ions by the reconnection (Chapter 5). The detailed background, assumptions, advantages and limitations of the two methods are presented in the following two subsections.

2.1. Large Scale Kinetic Simulation Used in the Study of Electron Acceleration in the Magnetotail

The LSK technique is a particle tracing simulation, which has been used in many studies [e.g., Ashour-Abdalla et al., 1993, 1994, 2011; Birn et al., 1998; Schriver et al., 1998, 2011]. In this technique, large numbers of particles are launched in time-varying electromagnetic fields and each of them follows single particle motion without causing any effect on the fields or other particles. As a result, the evolution of the distribution functions or the acceleration processes of these particles can be studied either statistically by calculating ensemble averages of the properties of the particles or by tracking individual particles.

For the LSK calculations used in this dissertation, the time-varying electromagnetic field was provided by the UCLA global magnetohydrodynamic (MHD) model [e.g., Raeder et al., 1995, 1998, 2001; El-Alaoui, 2001; El-Alaoui et al., 2009]. The input to the MHD simulation is a time series of the upstream solarwind conditions observed by spacecraft. The model
provides the magnetosphere wide response of the solar wind, magnetosphere and ionosphere system. The simulation is reasonably good at reproducing spacecraft observations in the magnetotail plasma sheet during substorms [e.g., El-Alaoui et al., 2009]. Detailed information about the MHD model including the boundary conditions on the ionosphere and the setup of numerical resistivity for triggering reconnection is discussed in El-Alaoui et al. [2009] and the references therein.

In Chapter 3 of this dissertation, the LSK simulation is used to study the electron acceleration near the X-line in the magnetotail. During the passage of the electrons in the tail, they can experience a variety of field configurations, including the reconnection region, thin current sheets, bursty bulk flow, and dipolarization fronts. Their motion can be described by guiding-center motion or full particle motion depending on the comparison between their gyro-radii and the local radius of curvature of the magnetic field. In the guiding-enter approximation charged particle motion in the electromagnetic field is approximated by the motion of the guiding center after removing the gyro-motion information. The equations giving the guiding-center motion are [Northrop, 1963]:

\[
\vec{v}_{gc,\perp} = \frac{\vec{E} \times \vec{B}}{B^2} + \frac{\mu}{q} \vec{B} \times \nabla B + \frac{p^2_{gc,\parallel}}{q m} \frac{\vec{B} \times (\vec{B} \cdot \nabla) \vec{B}}{B^4} ,
\]

(2-1-1)

\[
\frac{dp_{gc,\parallel}}{dt} = qE_{\parallel} - \mu \nabla_{\parallel} B ,
\]

(2-1-2)

where \( q \) and \( m \) are the charge and the mass, the subscripts “\( || \)” and “\( \perp \)” indicate the directions parallel and perpendicular to the local magnetic field respectively, \( \vec{v}_{gc,\perp} \) and \( p_{gc,\parallel} \) are the perpendicular velocity and the parallel momentum of the guiding center, \( \mu = \frac{mv_{\perp}^2}{2B} \) is the magnetic moment. Eq. (2-1-1) ignores the higher order terms. In the simulation, a fourth order
Runge-Kutta scheme is used for the time evolution. The time step is selected to ensure that the distance traveled by the guiding center in each time step is smaller than the spatial scale of the magnetic field gradient in any direction in the region of interest. For the full particle motion, the Lorentz-force equation is used:

\[
m \frac{d\vec{v}}{dt} = q\vec{E} + q(\vec{v} \times \vec{B}),
\]

where \(\vec{v}\) is the velocity of the charged particle. The time step is selected to be small enough to resolve the gyro-motion, e.g., \(\Omega \Delta t < 0.1\), where \(\Omega\) is the gyro-frequency.

When and where to switch between the guiding center motion and the full particle motion is determined by the adiabaticity parameter \(\kappa\) [Büchner and Zelenyi, 1986, 1989]. The adiabaticity parameter \(\kappa\) is defined as the square root of the ratio of the radius of curvature to the particle gyro-radius, i.e., \(\kappa = \sqrt{\frac{R_c}{\rho}}\). For \(\kappa >> 1\), the particle motion is adiabatic; for \(\kappa < 1\), the particle motion is non-adiabatic [Büchner and Zelenyi, 1986; Delcourt et al., 1996]; for \(1 < \kappa < 5\), the particle motion is not well described by the guiding-center method and is affected by many parameters [e.g., Delcourt et al., 1996]. For electrons, \(\kappa > 50\) in most of the regions in the magnetotail, while near the X-line, \(\kappa\) is frequently less than 10 and the electrons become non-adiabatic [Schriver et al., 1998]. The guiding-center calculation is only valid for adiabatic motion with large \(\kappa\). The full particle calculation is accurate and efficient when it is used only in the low \(\kappa\) range. In the simulation, the \(\kappa\) value at which the code switches between the guiding center and the full particle calculation is around 20-30. When the full particle information is transferred to the guiding-center scheme, the gyro-phase information is removed. When the particle is switched from the guiding center calculation to full particle motion, a random gyro-phase from 0° to 360° is added to the particle. To test the rationality of adding a
random phase, an experiment with thousands of particles was carried out which shows that the random gyro-phase fully covers the full 360° range [Schriver et al., 2011].

There are two ways in LSK simulations to study particle dynamics: (i) tracing the particles of interest individually and (ii) examining large groups of particles hitting virtual detectors in the region of interest. Both (i) and (ii) are frequently used in the study. The first approach can be used to study the energization history of the particles, but it has to be limited to a few particles due to the large amount of data output for every particle. The second approach is useful to compare with the satellite observations and confirm the simulation. The data from approach two is relatively small even though there are a large number of particles (e.g., ~10⁴), because the data are output only when the particles hit the detectors.

For the LSK simulations in this dissertation, the particles are launched impulsively near the X-line in the magnetotail plasma sheet. The virtual detectors are set earthward to the X-line including the locations with satellite data during a real event. The distribution function at the detector can be calculated based on the Liouville's theorem [e.g., Ashour-Abdalla et al., 1993, 1994; Richard et al., 1994, 2009]:

$$f(\nu) = \frac{(nv)_{MHD}}{\Delta v^3} \sum_{\nu_n} \frac{N}{N_{launch}} \frac{1}{A_l \Delta t_L} \frac{A_D}{A_D \Delta t_D},$$ (2-1-4)

where \((nv)_{MHD}\) is the number flux in the MHD, \(A_l\) and \(\Delta t_L\) are the launch area and the launch period, \(A_D\) and \(\Delta t_D\) are the detecting area and the detecting period, \(N_{launch}\) is the number of particles launched every period, \(\sum N\) is the number of particles detected every period, and \(\nu_n\) is the normal component of the velocity when a particle hits the surface of the virtual detector.
2.2. Implicit Particle-In-Cell (PIC) Simulation Used in the Study of Reconnection in the Presence of O\(^+\) Ions

One goal of this dissertation is to investigate the interaction between O\(^+\) ions and magnetotail reconnection, including the O\(^+\) effects on the reconnection and the O\(^+\) acceleration by the reconnection. This requires a self-consistent model to resolve the plasma kinetic behavior during the reconnection process and its feedback to the reconnection. To this end, the Particle-in-Cell (PIC) scheme [e.g., *Birdsall and Langdon*, 1985] is a good option. Due to the large separations (up to the order of 10\(^3\)) of masses of the electrons, protons, and O\(^+\) ions, the reconnection should be considered as a multi-scale plasma process involving the large separation of their spatial and temporal scales (up to the order of 10\(^3\) in the sense of gyro-motion). An explicit PIC scheme requires us to resolve the fastest electron response in the system, even though the focus point is on the long-term evolution of the reconnection and long-term behavior of the O\(^+\) ions. This can result in massive computation costs. Therefore, an implicit PIC scheme is applied to this work. As will be discussed later, an implicit scheme removes the constraint of the explicit scheme.

This section aims to introduce the general PIC scheme [e.g., *Birdsall and Langdon*, 1985] and the implicit PIC scheme [*Markidis et al.*, 2010]. It is divided into sub-sections, including the governing equations, the PIC method, the implicit scheme, and the numerical stability constraints.

2.2.1. Governing Equations

Since the plasma in the magnetotail plasma sheet is collisionless, it can be described by the Vlasov equation:
\[ \frac{\partial}{\partial t} f_s + \mathbf{v} \cdot \nabla f_s + \frac{q_s}{m_s} \left( \mathbf{E} + \frac{\mathbf{v} \times \mathbf{B}}{c} \right) \cdot \frac{\partial}{\partial \mathbf{v}} f_s = 0, \quad (2-2-1) \]

where the subscript “s” indicates the species, \( f_s = f_s(\mathbf{x}, \mathbf{v}, t) \) is the distribution function, \( q_s \) and \( m_s \) are the charge and the mass, \( \mathbf{E} = \mathbf{E}(\mathbf{x}, t) \) and \( \mathbf{B} = \mathbf{B}(\mathbf{x}, t) \) are the electric and magnetic fields. In this section, Gaussian units are used, which is reconciled to the iPic3D code. Besides the Vlasov equation, Maxwell’s equations are used to solve the evolution of the electromagnetic field. In iPic3D, Maxwell’s equations include terms to second order to guarantee that the divergence and curl equations are solved at the same time, which ensures both the uniqueness of the solution and the accuracy of the numerical solutions [Jiang et al., 1996; Ricci et al., 2002]:

\[ \nabla^2 \mathbf{E} - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{E} = \frac{4\pi}{c^2} \frac{\partial}{\partial t} \mathbf{J} + 4\pi \nabla \rho, \quad (2-2-2) \]

where the current density is \( \mathbf{J}(\mathbf{x}, t) = \sum_s \int q_s \mathbf{v} f_s(\mathbf{x}, \mathbf{v}, t) d\mathbf{v} \) and the charge density is \( \rho(\mathbf{x}, t) = \sum_s \int q_s f_s(\mathbf{x}, \mathbf{v}, t) d\mathbf{v} \). Once the electric field is solved, the magnetic field can be solved by considering the Faraday’s law \( \frac{\partial}{\partial t} \mathbf{B} = -c \nabla \times \mathbf{E} \). If the solution of the electromagnetic field initially satisfied the divergence equations \( \nabla \cdot \mathbf{E} = 4\pi \rho \) and \( \nabla \cdot \mathbf{B} = 0 \), then the two divergence equations are satisfied over time during the simulation.

2.2.2. Particle-in-Cell method

In the PIC method the Vlasov-Maxwell system is solved. In this method, the fields and moments are described on discretized grids, while the charged particles can move continuously in space. The computational cycle can be simply described as follows: (i) the fields and moments on the grids can be solved by using Maxwell’s equations; (ii) the electromagnetic
field on the grids then exerts a force on the charged particles and moves them by following
Newton’s second law; (iii) the charged particles in their new location form new charge and
current distributions and the new moments are converted to the grids to solve Maxwell’s
equations in (i) for the next cycle. The key issues in this method are how to define the “particles”
and the “communication” between the particles and the grids.

In a physical system, the number of particles is always huge and it is impractical to
simulate them directly even on the most powerful supercomputer in the world. In the tail plasma
sheet, the number density is \( \sim 0.2 \text{ cm}^{-3} \) for each species and the thickness of a thin current sheet
is in the order of \( \sim 10^3 \text{ km} \). The total number of one species is approximately \( \sim 10^{23} \), while one
of the recent reconnection studies using high performance supercomputer was only capable to
calculate the \( 10^{12} \) computation particles [Daughton et al., 2011]. In the PIC method, the concept
of computational particle has to be introduced. A computation particle is a small element in
phase space. It has finite size in space and a localized velocity. The distribution function can be
defined as:

\[
f_x(\vec{x}, \vec{v}, t) = \sum_{p=1}^{N_s} S(x - x_p)S(y - y_p)S(z - z_p)\delta(\vec{v} - \vec{v}_p),
\]

where, the subscript “p” is the label of computational particles with total number \( N_s \),
\( x_p = x_p(t) \),
\( y_p = y_p(t) \), \( z_p = z_p(t) \), \( \vec{v}_p = \vec{v}_p(t) \), \( S = S(x) \) is a shape function, \( \delta = \delta(\vec{v}) \) is Dirac’s delta
function. The shape function is a weight function that is used to interpolate the continuous
values to the distribution function based on the discretized distribution of the computational
particles with finite number. Because the number of the computational particles is much small
than the number of particles in a physical system, the shape function can effectively reduce the
noise caused by the rarefied distribution of computational particles in phase space. It can also be used to calculate the average electromagnetic field on the particles:

\[ \vec{E}_p = \int \vec{E}(\vec{x}) S(\vec{x} - \vec{x}_p) d\vec{x} \quad \text{and} \quad \vec{B}_p = \int \vec{B}(\vec{x}) S(\vec{x} - \vec{x}_p) d\vec{x}, \]  
(2-2-4)

Considering only the contribution of the electromagnetic field on the grids \((\vec{E}_g, \vec{B}_g)\), the average field on the particles can be described as:

\[ \vec{E}_p = \sum_g \vec{E}_g W(\vec{x}_g - \vec{x}_p) \quad \text{and} \quad \vec{B}_p = \sum_g \vec{B}_g W(\vec{x}_g - \vec{x}_p), \]  
(2-2-5)

where the weight function \(W(x_g - x_p)\) can be calculated by integrating the shape function in the cell in which the particle is located. It defines the communication between the particle and the surrounding grids. Also, the moments on the grids can be defined:

\[ \{ \rho_g, \vec{J}_g, \vec{I}_g \} = \sum_s \sum_p N_s q_s \{ \vec{1}, \vec{v}, \vec{v} \cdot \vec{v} \} W(\vec{x}_g - \vec{x}_p), \]  
(2-2-6)

where \(\rho_g\), \(\vec{J}_g\), and \(\vec{I}_g\) are the zero, first, and second order moments. Once the electromagnetic field \((\vec{E}_g, \vec{B}_g)\) on the particles is calculated, by substituting the distribution function of the computational particles (Eq. 2-2-3) to the Vlasov equation (Eq. 2-2-1), one can obtain the equations of motion:

\[ \frac{d\vec{x}_p}{dt} = \vec{v}_p \quad \text{and} \quad \frac{d\vec{v}_p}{dt} = \frac{q_s}{m_s} \left( \vec{E}_p + \frac{\vec{v}_p \times \vec{B}_p}{c} \right), \]  
(2-2-7)

where \(\vec{E}_p\) and \(\vec{B}_p\) are defined in Eq. (2-2-5).

### 2.2.3. Implicit Particle-in-Cell Scheme

The goal of the implicit PIC scheme is to eliminating the numerical stability constraints in the explicit scheme (which will be discussed in next subsection) without losing the fully kinetic
information. In the following, a brief review of the implicit PIC method is present based on the
Discussions in [Mason, 1981; Brackbill and Forslund, 1982; Vu and Brackbill, 1992; Ricci et
al., 2002; Lapenta et al., 2006; Markidis et al., 2010].

We start with the time discretization with a time step $\Delta t$. The current cycle is $t = n\Delta t$ and
the next cycle is $t + \Delta t = (n + 1)\Delta t$. For the variables used later, the superscripts “n” and “n+1”
indicate the current cycle and the next cycle respectively.

The equations of motion (Eq. 2-2-7) can be described as:

$$\bar{x}_p^{n+1} = \bar{x}_p^n + \bar{v}_p^{n+1/2} \Delta t$$  \hspace{1cm} (2-2-8)

$$\bar{v}_p^{n+1} = \bar{v}_p^n + \frac{q \Delta t}{m_s} \left[ \tilde{E}_p^{n+1}(\bar{x}_p^{n+1/2}) + \frac{\bar{v}_p^n}{c} \times \bar{B}_p^{n+1/2} \right]$$  \hspace{1cm} (2-2-9)

where $\bar{x}_p^{n+1/2} = (\bar{x}_p^{n+1} + \bar{x}_p^n)/2$, $\bar{v}_p^{n+1/2} = (\bar{v}_p^{n+1} + \bar{v}_p^n)/2$. The $\bar{v}_p^{n+1/2}$ can be expressed as a function
of $\bar{E}_p^{n+1}$:

$$\bar{v}_p^{n+1/2} = \tilde{\alpha}_s^n \cdot \bar{v}_p^n + \beta_s \tilde{\alpha}_s^n \cdot \bar{E}_p^{n+1}(\bar{x}_p^{n+1/2})$$  \hspace{1cm} (2-2-10)

$$\tilde{\alpha}_s^n = \frac{1}{1 + (\beta_s \bar{B}^n)^2} \left( \tilde{I} - \beta_s \tilde{I} \times \bar{B}_n + \beta_s^2 \tilde{I} \bar{B}^n \bar{B}_n \right)$$ and $\beta_s = \frac{q \Delta t}{2m_c}$  \hspace{1cm} (2-2-11)

where $\tilde{I}$ is the identity dyadic. The $\tilde{\alpha}_s^n$ includes the effect of the gyro-motion due to the
magnetic field.

For the electric field $\bar{E}_p^{n+1}$, it can be calculated from the time discretized Maxwell’s
equations:

$$\nabla \times \bar{E}_p^{n+1} + \frac{1}{c} \frac{\bar{B}_p^{n+1} - \bar{B}_n}{\Delta t} = 0$$  \hspace{1cm} (2-2-12)

$$\nabla \times \bar{B}_p^{n+1} = \frac{1}{c} \frac{\bar{E}_p^{n+1} - \bar{E}_n}{\Delta t} - \frac{4\pi}{c} \bar{j}_p^{n+1/2} = 0$$  \hspace{1cm} (2-2-13)
\[ \nabla \cdot \vec{E}^{n+1} = 4\pi \rho^{n+1} \quad (2-2-14) \]

\[ \nabla \cdot \vec{B}^n = \nabla \cdot \vec{B}^{n+1} = 0 \quad (2-2-15) \]

An exact implicit solution can be obtained by combining the Maxwell’s equations, the equations of motion and the definitions of the moments. However, it costs too many computational resources. In order to obtain a practical implicit method, an approximation has to be made. By combining the Maxwell’s equations, one can obtain:

\[ \vec{E}^{n+1} - (c\Delta t)^2 \nabla^2 \vec{E}^{n+1} = \vec{E}^n + (c\Delta t) \left( \nabla \times \vec{B}^n - \frac{4\pi}{c} \vec{J}^{n+1/2} \right) - 4\pi (c\Delta t)^2 \nabla \rho^{n+1} \quad (2-2-16) \]

The \( \rho^{n+1} \) and \( \vec{J}^{n+1/2} \) are the quantities that need to be solved to get \( \vec{E}^{n+1} \). For the moments in the continuous space, we have \( \{ \rho, \vec{J}, \vec{\Pi} \} = \sum_s \sum_p q_s \{ \vec{1}, \vec{v}, \vec{v}\vec{v} \} W(\vec{x} - \vec{x}_{sp}) \). For the \( \rho^{n+1} \), using Taylor expansion, one can obtain:

\[ \rho^{n+1} = \sum_s \sum_p q_s W(\vec{x} - \vec{x}_{sp}^{n+1}) \]

\[ = \sum_s \sum_p q_s [W(\vec{x} - \vec{x}_{sp}^n) + (\vec{x}_{sp}^n - \vec{x}_{sp}^{n+1}) \cdot \nabla W(\vec{x} - \vec{x}_{sp}^n)] + \frac{1}{2} (\vec{x}_{sp}^n - \vec{x}_{sp}^{n+1}) : \nabla \nabla W(\vec{x} - \vec{x}_{sp}^n) + \ldots \]

\[ = \sum_s \sum_p q_s [W(\vec{x} - \vec{x}_{sp}^n) - \Delta t \vec{v}_{sp}^{n+1/2} \cdot \nabla W(\vec{x} - \vec{x}_{sp}^n)] + \frac{(\Delta t)^2}{2} \vec{v}_{sp}^{n+1/2} \cdot \nabla \nabla W(\vec{x} - \vec{x}_{sp}^n) + \ldots \]

\[ \approx \rho^n - \Delta t \nabla \cdot \vec{J}^{n+1/2} \quad (2-2-17) \]

The last equation ignores the terms with higher order than \( \Delta t \). Similarly, \( \vec{J}_s^{n+1/2} \) can be obtained:

\[ \vec{J}_s^{n+1/2} \approx \vec{J}_s^n + \Delta t \frac{q_s \rho_s^n}{2m_s c} \vec{a}_s \cdot \vec{E}^{n+1} - \frac{\Delta t}{2} \nabla \cdot \vec{\Pi}_s^n \quad (2-2-18) \]
where \( \hat{J}^n_s = \sum_{p} q_s \hat{v}^n_{sp} W(\vec{x} - \vec{x}^n_{sp}) , \quad \hat{\Pi}^n_s = \sum_{p} q_s \hat{v}^n_{sp} \hat{v}^n_{sp} W(\vec{x} - \vec{x}^n_{sp}) \), and \( \hat{v}_{sp} = \vec{\alpha}_s \cdot \vec{v}^n_{sp} \). Note that (2-2-17) and (2-2-18) ignores the higher order terms of \( \Delta t \). The Taylor expansion was on the weight function and thus the shape function. Therefore, this method is only valid when the deviation in space caused by the high order terms is smaller than the spatial scale of the computational particle; otherwise, the Taylor expansion does not converge. By substituting (2-2-17) and (2-2-18) to (2-2-16), one can obtain

\[
(\vec{I} + \vec{\alpha}^n) \cdot \vec{E}^{n+1} - (c\Delta t)^2 \left[ \nabla^2 \vec{E}^{n+1} + \nabla \nabla \cdot (\vec{\alpha}^n \cdot \vec{E}^{n+1}) \right] = \vec{E}^n + (c\Delta t \left( \nabla \times \vec{B}^n - \frac{4\pi}{c} \hat{J}^n \right) - 4\pi (c\Delta t)^2 \nabla \hat{\rho}^n
\]

(2-2-19)

where \( \vec{\alpha}^n = \sum_s \frac{2\pi q_s \rho_s^n (\Delta t)^2 \vec{\alpha}^n}{m_s c} = \sum_s \frac{1}{2} \left( \omega_{ps} \Delta t \right)^2 \vec{\alpha}^n \), \( \omega_{ps} = \sqrt{\frac{4\pi q_s \rho_s^n}{m_s c}} \) is the plasma frequency for species “s”, \( \hat{\rho}^n = \rho^n - \Delta t \nabla \cdot \hat{J}^n \). In the Eq. (2-2-19), \( \vec{E}^{n+1} \) is the only variable unknown. This equation together with (2-2-12) can be solved by discretizing in space to obtain the electromagnetic field on the grids at \( (n+1)\Delta t \).

Although the solution of the second order Maxwell’s equations with proper initial and boundary conditions theoretically holds for the divergence equations \( \nabla \cdot \vec{E} = 4\pi \rho \) and \( \nabla \cdot \vec{B} = 0 \), the microscopic inconsistencies between \( \rho \) and \( \vec{J} \) due to communication between the particles and the grids can lead to small violations of the divergence equations. A correction to the electric field can be made by introducing a potential \( \phi \) \( [\text{Boris, 1970; Birdsall and Langdon, 1985}], \) which gives the correct electric field \( \vec{E}' \):

\[
\vec{E}' = \vec{E} - \nabla \phi
\]

(2-2-20)

Based on the Gauss’ law, the potential \( \phi \) can be obtained by solving:
\[ \nabla^2 \Phi = \nabla \cdot \vec{E} - 4\pi \rho \]  
(2-2-21)

### 2.2.4. Numerical Stability Constraints

As mentioned earlier, the goal of developing the implicit PIC scheme is to remove the restrictive constraints on the numerical stability necessary for the explicit PIC scheme.

In the explicit PIC scheme, there are three stability constraints:

(i) The time step \( \Delta t \) should be small enough to resolve the fastest propagation of the oscillating signal between the grid spacing \( \Delta x \):

\[
\Delta t < \frac{\Delta x}{v_{\text{max}}} = \frac{\Delta x}{c},
\]

where \( c \) is the speed of light, the fastest wave speed in the plasma system. This is the Courant–Friedrichs–Lewy (CFL) condition [Courant et al., 1928].

(ii) To avoid the aliasing due to the finite time step, a constraint is introduced [Birdsall and Langdon, 1985]:

\[
\omega_{\text{ps, max}} \Delta t < 2 \Rightarrow \omega_{\text{pe}} \Delta t < 2,
\]

where \( \omega_{\text{ps, max}} \) is the maximum plasma frequency for all species, which is typically the electron plasma frequency \( \omega_{\text{pe}} \). This constraint can be understood as the requirement of resolving the fastest oscillation in the system.

(iii) The loss of information due to the interpolation between the particles and the grids results in a finite grid instability. A constraint is added to limit this instability:

\[
\Delta x < \zeta \lambda_{\text{De}},
\]

where \( \lambda_{\text{De}} \) is Debye length. The constant factor \( \zeta \) is the order of one. Its value depends on the specific scheme.
In the implicit PIC scheme, the constraints (2-2-22) and (2-2-23) for the explicit scheme are removed. As described in Brackbill and Forslund [1982], by solving the equations backward in time, the implicit scheme can effectively prevent the temporal growth of the numerical instability comparing with the forward solution. Replacing the explicit constraints, are two constraints for both the stability and accuracy in the implicit scheme:

(i) The Taylor expansions (2-2-17) and (2-2-18) of the weight function are truncated as an approximation. The requirement of the convergence of the expansion gives the constraint [Brackbill and Forslund, 1982]:

\[ \Delta t < \frac{\Delta x}{v_{th,e}}, \quad (2-2-25) \]

where the \( v_{th,e} \) is the electron thermal velocity. This constraint limits the motion of a particle so that it crossed no more than one cell in a time step.

(ii) The explicit constraint of the finite grid instability is replaced by [Brackbill and Forslund, 1982]:

\[ \frac{\Delta x}{\Delta t} < \varphi v_{th,e} \quad (2-2-26) \]

This allows larger grid spacing when the larger time step is used.

As a result, in the implicit scheme, the time step can be larger than the fastest oscillation period in the system. Once a time step is chosen, the fast processes occurring within the time step are kept but averaged. This property does not exist in the fluid scheme or hybrid scheme, in which such fast processes are entirely removed. The verification and validation of this implicit PIC scheme are discussed in Lapenta et al. [2006] and Markidis et al. [2010].

As mentioned in the Chapter 1, reconnection is a process involving multiple scales: from MHD scales to ion kinetic scales to electron kinetic scales. A proper reconnection study has to
consider both electron and ion kinetics. Although the explicit PIC scheme can address the reconnection problems reasonably, it is difficult to study long-time effects, e.g., the temporal scales of heavy ion behavior, while maintaining the constraints of resolving the electron scales, due to the huge computation cost. Based on the discussion above, in the implicit PIC scheme, when the time step is larger than the electron scales, the electron kinetic effects are still included, even though the electron scales are not fully resolved. Therefore, with appropriately large time steps, an implicit PIC scheme allows us to efficiently study the long-term effects, e.g., reconnection evolution and heavy ion behavior, without losing the electron kinetics. This is the reason why the implicit PIC scheme is chosen to address the problems about the interaction between the oxygen ions and the reconnection in this dissertation. The specific simulation set-up for studying the reconnection in the presence of the oxygen ions will be discussed later in Chapters 4 and 5.
Chapter 3

Electron Acceleration in Realistic Magnetotail X-line Configurations

3.1. Introduction

In this chapter, we investigate electron acceleration during two substorm events featuring different magnetotail configurations and particle acceleration processes. The study utilizes in situ observations, global MHD simulations, and large-scale-kinetic (LSK) simulations. The two substorms in question occurred on 15 February 2008 and 15 August 2001. The 15 February 2008 event was observed by the THEMIS spacecraft at \(X_{\text{GSM}} \sim -10\ \text{R}_E\), while the 15 August 2001 event was observed by Cluster spacecraft at \(X_{\text{GSM}} \sim -20\ \text{R}_E\). In the geocentric solar magnetospheric (GSM) coordinate system, in which the X axis points from the Earth to the Sun, the Z axis is the projection of the Earth's magnetic dipole to the plane perpendicular to the X axis and pointing in the geographic north direction and the Y axis is defined by the right-hand rule. All results in this dissertation use the GSM coordinates. The two events have different solar wind driving conditions: the 15 February 2008 event features an interplanetary magnetic field (IMF) \(B_z\) mostly between \(-2\) and \(2\ \text{nT}\) and a solar wind speed above \(600\ \text{km/s}\), while the 15 August 2001 event exhibits a relatively long interval of very weak southward IMF component before it reaches the minimum around \(-2\ \text{nT}\) and a solar wind speed around \(450\ \text{km/s}\). We used a global magnetohydrodynamic (MHD) simulation driven by observed solar wind conditions for each event and found that the two events had very different magnetotail configurations. During the 15 February 2008 event, a localized neutral line with a high speed (~300 km/s) and a narrow (<5 R_E in dawn-dusk direction) flow channel are found in the MHD simulation. A large cross-tail neutral line and low-speed (~100 km/s) flow across a broad region
of the tail (>15 \( R_E \) in dawn-dusk direction) existed during the 15 August 2001 event. We used the global three-dimensional time-varying electric and magnetic fields from the MHD simulation to carry out large-scale kinetic (LSK) simulations, in which thousands of electron trajectories were calculated. As a result, the spatial evolution of the energetic electron distribution functions in the near-Earth plasma sheet can be mapped. This allows us to understand under what magnetotail configurations different electron distribution functions are formed during each substorm event.

3.2. Observations

In this section we present observations from the THEMIS and Cluster spacecraft during the February 15, 2008 and August 15, 2001 substorms. We now present the observational details for these two events.

3.2.1. 15 February 2008 Event

On 15 February 2008 THEMIS satellites P3, P4 and P5 were located at \((-10.7, 1.6, -3.5)\), \((-9.8, 2.4, -3.5)\) and \((-9.0, 1.9, -3.1)\) \( R_E \) respectively. Multiple earthward propagating dipolarization fronts (DFs) were detected [Zhou et al., 2009; Ashour-Abdalla et al., 2011]. The AE index between 0300 and 0330 UT, provided by the World Data Center (Kyoto), showed an increase from 100 to 300 nT and it remained around 300 nT from 0330 to 0430 UT after which it started to decrease. THEMIS Data used in this chapter are measured by the following instruments: the Flux Gate Magnetometer (FGM) [Auster et al., 2008], electron distribution functions in the energy range from 5 eV to 30 keV from the Electrostatic Analyzer (ESA) [McFadden et al., 2008], and electron differential energy flux with high energies from 30 to 421 keV from the Solid State Telescope (SST) [Angelopoulos, 2008].
We provide an overview of this event observed by the THEMIS P4 satellite in Figure 3-1. Figure 3-1a, 3-1b and 3-1c show the magnetic field components in GSM coordinates, differential energy flux of energetic electrons (> 30 keV) from SST and differential energy flux of electrons (< 30 keV) from ESA, respectively. Because P3 and P5 were close to P4 and detected similar results, we show here only the P4 data. THEMIS P4 detected two dipolarization fronts, characterized by rapid increases in $B_z$, at 0357:00 and 0357:38 UT. Behind the DFs, electron fluxes at energies larger than 30 keV quickly increase and at energies less than 30 keV decrease. Zhou et al. [2009] have discussed this event in detail and showed that lower hybrid drift (LHD) waves and electron cyclotron harmonic (ECH) waves were associated with the DFs and are candidates for electron energization. By comparing electron flux between simulations and observations, Ashour-Abdalla et al. [2011] analyzed the acceleration mechanism of energetic electrons related to the DFs for this event and concluded that electrons were initially accelerated near the x-line up to several keV energies and additional energization up to ~ 100
keV occurred due to betatron acceleration as the electrons moved earthward with the DF. Here, the THEMIS P4 observations shown in Figure 3-1 will be compared with global kinetic simulations in order to further understand physical processes occurring during the substorm and to compare with the 15 August 2001 substorm event.

### 3.2.2. 15 August 2001 Event

During the event on 15 August 2001, the four Cluster spacecraft were located near the post-midnight plasma sheet in the magnetotail \( \sim 18 \, R_E \) from the Earth. Multiple DFs were observed by the four Cluster spacecraft between 0800 and 0900 UT. These DFs were associated with a well-documented substorm event [McPherron et al., 2002; Kivelson et al., 2005; Weygand et al., 2005]. The expansion phase onset for this event was just before 0730 UT and the Cluster observations corresponded to an expansion of the auroral bulge to higher latitudes [Hwang et al., 2011]. The AE index, provided by the World Data Center (Kyoto), increased to \( \sim 200 \, \text{nT} \) at around 0810 UT, and then it fluctuated around 200 nT with three obvious peaks in the interval between 0810 and 0900 UT. Data are provided by the following instruments onboard Cluster: magnetic field from the Flux Gate Magnetometer (FGM) [Balogh et al., 2001], electron differential energy flux and distribution functions with energy range from 10 eV to 26.5 keV from the Plasma Electron and Current Experiment (PEACE) [Johnstone et al., 1997], and electron differential energy flux with energies from 39.2 to 244 keV from the Research with Adaptive Particle Imaging Detectors (RAPID) [Wilken et al., 2001].
Figure 3-2 Overview of Cluster 1 observations for the 15 August 2001 event: (a) magnetic field components (b) electron pitch angle distribution from the PEACE instrument (c) differential energy flux spectrum from the PEACE instrument (d) electron differential energy flux from the RAPID instrument. The magenta dashed lines indicates the six DFs (after Hwang et al., 2011).

An overview of this event observed by the Cluster 1 (C1) satellite is shown in Figure 3-2. Figure 3-2a, b, c and d, show the magnetic field components in GSM coordinates, electron pitch angle distribution from PEACE, differential energy flux spectrum of electrons with low energy (< 26 keV) from PEACE and that of energetic electrons (> 26 keV) from RAPID, respectively. The C1 spacecraft was located at (−18.0, −5.4, −1.0) $R_E$ GSM. Since the other three Cluster spacecraft obtained similar results, we show only the C1 data here. As shown in Figure 3-2a, from 0817 to 0841 UT, six dipolarization fronts (DFs) were observed. In Figure 3-2b, after the
1st, 2nd, 3rd and 5th DF, bidirectional streams were observed. In Figure 3-2c and 3-2d, differential energy flux with energies less than 2 keV increases after the 1st, 3rd and 5th DF, while differential energy flux with energies between 2 and 94.5 keV decreases after the 1st, 2nd, 3rd and 5th DF. The features of the multiple DFs during this event are discussed in detail by Hwang et al. [2011]. In this chapter, we focus on the magnetotail configurations and associated non-adiabatic effects on electrons during this event, which have not been discussed before. In section 3.4, the C1 observations are compared to the MHD and LSK simulations.

3.2.3. Solar Wind Conditions for the Two Events

The events selected for this study have different solar wind conditions which resulted in different magnetotail configurations. Figure 3-3 shows the solar wind conditions, including interplanetary magnetic field (IMF) components in GSM (Geocentric Solar Magnetospheric) coordinates, ion velocity components in GSM coordinates, and density, for the two events. For the 15 February 2008 event, solar wind conditions were detected by Cluster spacecraft: magnetic field components are from FGM, and ion bulk velocity and density from the Hot Ion Analyser (HIA) of Cluster Ion Spectrometry (CIS) [Rème et al., 2001]. For the 15 August 2001 event, solar wind conditions were provided by the Advanced Composition Explorer (ACE) satellite: magnetic field components from the Magnetic Fields Experiment (MAG) [Smith et al., 1998], solar wind velocity and density from the Solar Wind Electron Proton Alpha Monitor (SWEPAM) [McComas et al., 1998]. These observations have been propagated to the upstream sunward boundary of the MHD simulation box at 20 \( R_E \). The intervals of interest are shown in blue. During the 15 February 2008 event, the IMF \( B_x \) is dominant before the \( B_y \) component becomes dominant at around 0313; the IMF \( B_z \) ranges from about \( -2 \) to 2 nT and the solar wind speed is above 600 km/s. During the 15 August 2001 event, the IMF \( B_z \) also ranges from \( -2 \) to
2 nT, but a dominant IMF $B_y$ component exists over a relatively much longer interval and the solar wind speed is around 450 km/s. The solar wind conditions for each event are used as input to drive the MHD simulations, which are discussed in next section.
Figure 3-3 Solar wind observations from Cluster on 15 February 2008 (top set of panels) and ACE on 15 August 2001 (bottom set of panels). For each event, the top three panels give the magnetic field in GSM coordinates. The next three panels show the flow velocity and the bottom panel gives the density. These observations have been propagated to the upstream sunward boundary of the MHD simulation box at 20 $R_E$.

3.3. Simulations

Global three-dimensional kinetic simulations that combine the magnetohydrodynamic (MHD) and large scale kinetic (LSK) techniques have been carried out for each event. Details of the MHD and LSK simulation are now discussed.
3.3.1. MHD Simulation

We have used the University of California, Los Angeles (UCLA) global MHD code to simulate the magnetospheric response to solar wind conditions. This model has been discussed in detail previously [e.g., Raeder et al., 1995, 1998, 2001; El-Alaoui, 2001, El-Alaoui et al., 2009], so we provide only a brief discussion here. This model is a three-dimensional single-fluid code and uses observed solar wind conditions to drive the simulation. The simulation system size in GSM coordinates is from 20 to −300 $R_E$ in $x$, and from −55 to 55 $R_E$ in $y$ and $z$. The numbers of grids in the three dimensions are $400 \times 240 \times 240$ for 15 February 2008 event and $416 \times 300 \times 300$ for 15 August 2001 event. The minimum grid dimensions are $0.20 \times 0.20 \times 0.20$ $R_E$ for the 15 February 2008 event and $0.20 \times 0.15 \times 0.15$ $R_E$ for 15 August 2001 event. The inner boundary is a spherical shell at 2.7 $R_E$, representing the Earth’s ionosphere, coupled with the magnetospheric domain through an ionospheric potential equation [Raeder et al., 1995]. By solving resistive MHD equations as an initial value problem, the response of the magnetosphere, including the time evolution of the electric and magnetic fields, flow, current density and number density, can be calculated. In order to model the effects of current-driven instabilities, a resistive term together with a convective term is included in Ohm’s law, i.e., $E = -V \times B + \eta j$, where $V$, $B$ and $j$ are flow velocity, magnetic field and current density respectively. The anomalous resistivity $\eta$ is defined as $\eta = \alpha j''$, where $j'' = \frac{|\Delta|}{|B|+\epsilon}$. If $j' < \delta$, $j' = 0$. The $\delta$ is determined by limiting the $\eta \neq 0$ region at very few grids with very strong current sheets. Similar expressions for anomalous resistivity have been used in previous magnetic reconnection studies [e.g., Sato and Hayashi, 1979; Hoshino, 1991]. Solar wind observations shown in Figure 3 are used as input to drive the MHD simulations.
3.3.2. LSK Simulation

In order to determine electron kinetic particle effects, we use LSK in which thousands of electrons are launched as test particles and their trajectories are calculated in the time-dependent electric and magnetic fields obtained from the MHD simulations for each event. This method has been discussed in detail in previous studies [e.g., Ashour-Abdalla et al., 1993, 2011; Schriver et al., 1998]. The electron trajectories are calculated by using a combination of guiding center and full particle orbit calculations [Schriver et al., 2011]. The adiabaticity parameter $\kappa$, defined as the square root of the local magnetic field radius of curvature divided by the local gyroradius of the electron [Büchner and Zelenyi, 1986], is used to determine when and where to switch between guiding center and full particle orbit calculations. In the simulations here, when $\kappa < 10$, the full particle orbit calculation is used; otherwise, the guiding center equations are used. Virtual detectors are set up at the locations of interest in the simulation. Once the electrons hit any of the virtual detectors, electron information, including time, position, parallel and perpendicular velocity components (relative to background magnetic field), kinetic energy and $\kappa$ parameter, are collected in order to calculate the differential energy flux and distribution function.

For both simulations the launch regions for the test electrons were near the plasma sheet. They were determined by tracing particle trajectories backward in time from spacecraft locations and by considering the location of the reconnection region and flow channels in the MHD simulations. The launch sites were downstream from the neutral line in a region with small adiabaticity parameter $\kappa$.

For the 15 February 2008 event, the launch region was a square with $-19 \ R_E < x < -21 \ R_E$ and $1 \ R_E < y < 4.5 \ R_E$ centered on the current sheet earthward from the neutral lines. The center
of the current sheet was defined as the surface of maximum plasma thermal pressure in the MHD simulation [Ashour-Abdalla et al., 2002]. The location of the near Earth neutral line for this event is approximately between \(x \sim -25\) and \(-30\) \(R_E\). We launched 65,000 particles every 20 seconds. The electrons had a Maxwellian distribution with thermal energy \(E_{th} = 1\) keV. We launched a Maxwellian distribution without a high energy tail although such a tail is sometimes seen in observations. However, since for this event there were no spacecraft observations near the neutral line, we used a standard Maxwellian with a typically observed temperature of 1 keV in the plasma sheet for our launch distribution. As we will see later in this chapter, using such a launch distribution gives results that compare well with the observations.

For the 15 August 2001 event, the launch region again is a square with \(-9\) \(R_E < y < -1.5\) \(R_E\) and \(-6\) \(R_E < z < 0\) \(R_E\) at \(x = -22\) \(R_E\) in GSM coordinates. The location of the near Earth neutral line for this event is approximately from 20 to 25 \(R_E\). For this event we launched 70,000 particles every 20 seconds. It is worth noting that the near Earth plasma sheet for this event was very thin and with a large tilt angle (\(\sim 13^\circ\) to the dawn in GSM coordinates at \(x \sim -20\) \(R_E\)). Because of this configuration we launched the electrons in a rectangular region in the \(y\)-\(z\) plane covering the current sheet rather than on the surface of maximum thermal pressure, which results in adequate loading of the plasma sheet with electrons. The launched electrons had a Maxwellian distribution with thermal energy \(E_{th} = 4\) keV; this is based on observations from C1, which was close to the neutral line in the MHD simulation.
3.4. Results

3.4.1. Comparisons between the Observations and the Simulations

For the two events in the study, we launched electrons near the neutral line and wanted to see what happens when they propagate earthward. To validate the simulation results and carry out further analysis of energization mechanisms operating during the different substorms, we perform comparisons between the MHD and LSK simulation results and the spacecraft observations for both events. The similarities between the simulations and observations for the two events are summarized in Table 3-1. The details of these comparisons are discussed in the following subsections.

Table 3-1 List of similarities between observation and simulation for 15 February 2008 event and 15 August 2001 event

<table>
<thead>
<tr>
<th>15 February 2008 Event</th>
<th>15 August 2001 Event</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Observation</strong></td>
<td><strong>Simulation</strong></td>
</tr>
<tr>
<td>Two DFs 0357~0358 UT observed by THEMIS P4 (start at 0357:00, 0357:40 UT)</td>
<td>Two DFs 0356~0359 UT in MHD (start at 0356:40, 0358:20 UT)</td>
</tr>
<tr>
<td>Energetic electron flux ($&gt;26$ keV) decreases after the 1st, 2nd, 3rd and 5th DFs, whereas it</td>
<td></td>
</tr>
</tbody>
</table>
3.4.1.1. 15 February 2008 Event

The simulation results for this event are shown in Figure 3-4. In Figure 3-4a-f, the color coded contours of $B_z$ are shown on the surface of maximum pressure at different times. The locations of the P3, P4 and P5 spacecraft are indicated as black circles shown in Figure 3-4a-f and the flow vectors are indicated by white arrows. Two dipolarization structures successively engulf the three spacecraft. The first DF encounters P4 and P5, but not P3, and the second DF engulfs P3, P4 and P5 in order, consistent with the observations (Figure 3-4g).
Figure 3-4 Results from the MHD and LSK simulations for the 15 February 2008 event: (a)-(f) $B_z$ spectrograms on the surface of maximum pressure in the MHD simulation at different times. White arrows are flow vectors. Circles indicate THEMIS P3, P4 and P5 locations (after Ashour-Abdalla et al., 2011). (g) $B_z$ observations by P3, P4 and P5 spacecraft. Vertical dashed lines indicate the times of (a)-(f) spectrograms. (h) and (i) distribution functions $f(v_{\parallel})$ and $f(v_{\perp})$. The sizes of $v_{\parallel}$ and $v_{\perp}$ bins are $1 \times 10^4$ and $2 \times 10^4$ km/s respectively. See text for details.

The differential energy flux and distribution functions are calculated by analyzing LSK electron hits on the virtual detectors [Ashour-Abdalla et al., 1993; Richard et al., 2009]. For this event, comparisons of differential energy flux between the observations and simulations have been discussed in Ashour-Abdalla et al. [2011]. Here we compare distribution functions between the observations and simulations in Figure 3-4h and 3-4i. Colored solid lines are distribution functions with energies less than 30 keV observed by the ESA instrument onboard THEMIS P4. The times range from 30 seconds before to 30 seconds after the dipolarization front at 0357:00 UT. Rainbow colors indicate the time with blue lines before the DF and red lines are after the DF. In the LSK simulation, to calculate the distribution function accurately, a large number of electron counts is required, meaning we need a large collecting area and/or a long collecting time. The simulated distribution functions for this event are calculated by using electron information collected over a $2 \times 2 R_E$ region centered at $(x, y) = (-9.8, 2.4) R_E$ on the surface of maximum pressure and in the intervals of 30 seconds before and after the DF around 0357:00 UT, shown by dashed lines with blue and red colors in Figure 3-4h and 3-4i, indicating before and after the DF respectively. We did not launch a cold plasma component, so we cannot compare simulation results of the lower velocity part of the distribution ($v < 2 \times 10^4$ km/s) with
the observations. Both observed and simulated $f(v_{\parallel})$ show larger magnitude before (blue) the DF than that after (red) the DF and they show similar trends and magnitudes in the velocity range between $2 \times 10^4$ km/s and $1 \times 10^5$ km/s. The observed $f(v_{\perp})$ shows a larger magnitude before the DF than after the DF from $3 \times 10^4$ km/s to $6.5 \times 10^4$ km/s (from ~ 2.6 to 12 keV). Similarly, the simulated $f(v_{\perp})$ before the DF shows a larger magnitude than after the DF from $4 \times 10^4$ km/s to $9.5 \times 10^4$ km/s (from ~ 4.6 to 26 keV).

3.4.1.2. 15 August 2001 Event

For the second event (observations shown in Figure 3-2), six DFs were observed by Cluster 1 at 0821:30, 0825:00, 0828:30, 0832:30, 0835:00 and 0839:50 UT. Figure 3-5a shows MHD $B_z$ detected at (-18.0, -5.4, -1.0) $R_E$ in GSM, the C1 spacecraft location (black line) in the simulations, and Z at other locations. Red and blue lines indicate northward (+) and southward (-) shifts in Z respectively. The DFs were not very strong (~10 nT) in the observations, and they are even weaker (~4 nT) in the simulation due to the smoothing between finite grid points. Thus, we have shifted the z location to find clear DFs in the simulation. The $B_z$ at $z + 0.25$ $R_E$ and $z + 0.50$ $R_E$ in the MHD simulation shows six DFs at 0819:00, 0822:30, 0827:00, 0832:10, 0834:50 and 0839:40 UT.
Figure 3-5 Results from the MHD and LSK simulations for the 15 August 2001 event: (a) $B_z$ component obtained at C1 location and shifted-z locations in MHD simulation (b) Electron differential energy flux from LSK simulation with the same energy channels as RAPID (c) and (d) Distribution functions $f(v_{||})$ and $f(v_{\perp})$. The sizes of $v_{||}$ and $v_{\perp}$ bins are $1.1 \times 10^4$ and $1.6 \times 10^4$ km/s respectively. See text for details.

In Figure 3-5b, the simulated differential energy flux of electrons in high ($> 26$ keV) energy channels show flux dips at high energies from 39.2 keV to 68 keV behind the 1st, 2nd,
and 4th DF similar to the observations (Figure 3-2). This simulated particle result is calculated from spacecraft on the maximum pressure surface. Figure 3-5c and 3-5d show a comparison of distribution functions in low energies (< 26 keV) from the LSK simulations and the observations in a format similar to Figure 3-4 with colored solid lines showing observations with energies less than 26 keV from the PEACE instrument onboard C1. As in Figure 3-4, different colors indicate different detection times during the DF passage and the dashed lines are the LSK simulation results. Observations are shown at times around the first DF passage and the LSK results are shown around the 2nd DF passage. The second DF in the simulation was selected because of strong $B_z$ and electron signatures. Again, we cannot compare simulation results with the cold plasma peak in the observations since cold plasma was not launched in the LSK simulations. For the observations, both $f(v_{||})$ and $f(v_{\perp})$ show a larger magnitude before the DF than after the DF from $2 \times 10^4$ km/s to $1 \times 10^5$ km/s. Simulated $f(v_{||})$ also shows larger magnitude before the DF than after the DF in the range $2 \times 10^4$ km/s to $1 \times 10^5$ km/s, and the simulated $f(v_{\perp})$ before the DF shows larger magnitude than $f(v_{\perp})$ after the DF in the entire range up to $1 \times 10^5$ km/s. Both $f(v_{||})$ and $f(v_{\perp})$ in the simulation have similar magnitudes and slopes as those in the observations. It is worth noting that the fluxes in the lower energy channels decrease after the DF in both observations and the simulation, while those in the higher energy channels increase in the observations indicating acceleration. We did not compare the distribution functions in higher energy channels (> 30 keV) in the simulation with the observations because there were too few particles in these channels to make an accurate comparison. For example, in the observation of Figure 3-2d, the differential energy fluxes are generally around $10^4$ eV/(cm$^2$ s sr eV) in the energy range between 39.2 and 50.5 keV, $10^3$~$10^4$
eV/(cm² s sr eV) at 68 keV, and 10²–10³ eV/(cm² s sr eV) in the energy range from 94.5 to 244 keV. In the simulation as shown in Figure 3-5b, the differential energy fluxes are around 10⁴ eV/(cm² s sr eV) at 39.2 keV, 10³ eV/(cm² s sr eV) at 50.5 keV, and 10² eV/(cm² s sr eV) at 68.0 keV. There are insufficient electrons to calculate the flux in the energy channels higher than 68 keV.

When we compare the LSK results and the observations, two points need to be noted:

1. The electrons in the simulation generally follow the flows in the tail. The simulated flows are not accurate enough in space and time for detailed comparison. As a result, the simulated flux of electrons can be shifted in pace and time with respect to the observations.

2. The simulated electrons are launched in specific regions. The launch regions cannot cover all of the possible source regions in the observation. Thus not all of the particles that are observed will be included in the launches. Some will be from areas outside of our launch region. These two points are important for understanding the comparison between LSK results and the observations. For example, in the simulation of the 15 August 2001 event, there are strong particle signatures after the second DF but weak particle signatures after the third DF. In the simulation the flow just misses C1 with the result that too few particles reach the simulated spacecraft position.

3.4.2. Magnetotail Configurations

In order to further study the effects of the magnetotail configuration on electrons during the two different substorm events, a comparison of magnetotail parameters is shown in Figure 3-6. The plots are on the surface of maximum thermal pressure from the MHD simulations for 15 February 2008 on the left and 15 August 2001 on the right. The color contours give the $B_z$ component. Arrows and white solid lines indicate flows and neutral lines respectively. Typical
earthward flows in both events are highlighted by dashed blue lines. Spacecraft locations are indicated as small circles. For the 15 February 2008 event (left panels) the magnetotail configurations before the first DF (0356:00 UT), during the first DF (0357:00 UT), and after the two DFs (0358:00 UT) are shown. For the 15 August 2001 event (right panels) the magnetotail configurations before (0822:00 UT), during (0823:00 UT), and after the second DF (0824:00 UT) are shown. Comparing the two events shows that the magnetotail configurations are very different. For 15 February 2008, there are localized neutral lines and strong flows (> ~300 km/s) in a narrow flow channel (< 5 \(R_E\) in the \(y\) direction), while the 15 August 2001 event exhibits cross-tail neutral lines and slow flows (< 100 km/s) over a broad (> 15 \(R_E\)) region in the \(y\) direction. In addition, the 15 August 2001 event has a much thinner plasma sheet than the 15 February 2008 event (not shown here). Such different magnetotail configurations are mainly caused by the different solar wind driving conditions for the two events.
Figure 3-6 $B_z$ component on the surface of maximum pressure for 15 February 2008 event (left panels) and 15 August 2001 event (right panels) from the MHD simulation.
3.4.3. Electron Distribution Functions

To investigate the effect these differences have on electrons for the two events, we calculated distribution functions parallel and perpendicular to the local magnetic field from the LSK simulation results. The distribution functions are compared between the two events at $x$ locations between $-10 \, R_E$ and $-18 \, R_E$. Figure 3-7 shows the spatial evolution of parallel and perpendicular velocity distribution functions $f(v_\parallel)$ and $f(v_\perp)$ on the surface of maximum thermal pressure for the two events, with $f(v_\parallel)$ in red and $f(v_\perp)$ in blue. We used virtual detectors that completely spanned the dipolarization fronts. The $y$ range for the virtual detector is from 0 to 6 $R_E$ for 15 February 2008, and from $-8$ to 1 $R_E$ for 15 August 2001. The two intervals are from 0357:00 to 0358:00 UT for 15 February 2008 event and from 0836:00 to 0837:00 UT for 15 August 2001, when there were enough electrons covering the regions between $-10 \, R_E$ and $-18 \, R_E$. It is worth noting that in general, the distribution functions at the locations closer to Earth are formed by the electrons launched earlier in the run. The distribution functions for 15 February 2008 are anisotropic with $f(v_\parallel) > f(v_\perp)$ from $x = -10$ to $x = -16 \, R_E$. Such anisotropy could be formed by betatron acceleration [e.g., Ashour-Abdalla et al., 2011].

The distribution functions at $x = -17$ and $x = -18 \, R_E$ are generally isotropic. One possible reason for this is that these two locations are very close to the launch region and thus display the isotropic form of the launch distributions. For 15 August 2001, the distribution functions are anisotropic with $f(v_\parallel) > f(v_\perp)$ from $x = -10 \, R_E$ to $x = -15 \, R_E$, especially for the high energy tails, although they are isotropic at $x = -16 \, R_E$ and $f(v_\perp) > f(v_\parallel)$ at $-17 \, R_E$ and $-18 \, R_E$. This result is very different from the 15 February 2008 event.
Figure 3-7 Distribution functions on the surface of maximum pressure at different $x$ locations with the parallel distribution function shown in red and the perpendicular distribution function in blue.
3.4.4. Non-adiabatic Effects in the Magnetotail

In Figure 3-8a and 3-8b we evaluate whether the electrons could experience non-adiabatic acceleration during these two events. The figures show the color coded adiabaticity parameter $\kappa$ (defined in section 3.3.2) for 2-keV electrons on the surface of maximum pressure for the two events. The locations of the THEMIS P4 and Cluster 1 spacecraft are shown by small circles for both events. Higher $\kappa$ (in green, yellow and red) means the electrons are adiabatic, lower $\kappa$ (in blue) indicates the electrons can undergo non-adiabatic motion in the magnetotail current sheet. Non-adiabatic effects include acceleration by the magnetic reconnection electric field near neutral lines and scattering by small scale electric fields in the plasma sheet. For both events, electrons were launched in the region with small $\kappa$ (near the neutral lines) as indicated in Figure 3-8. Note that for the 15 February 2008 event, the neutral line is very dynamic and the launch region shown in Figure 3-8a is the region where the neutral line was located in earlier times (every 20 seconds after 0347:00 UT).
Adiabatic Parameter $\kappa$ on Maximum Pressure Surface at 03:57:00UT for 2keV Electrons

Feb 15, 2008

Adiabatic Parameter $\kappa$ on Maximum Pressure Surface at 08:23:00UT for 2keV Electrons

Aug 15, 2001
Figure 3-8 Adiabatic parameter κ on the surface of maximum pressure. White solid line represents the neutral line.

For 15 February 2008, in general, small κ regions mainly cover the tail region beyond 18 \( R_E \) (or slightly closer to earth after midnight). There is only a narrow (< 3 \( R_E \)) region with small κ earthward of the launch region and it ends at around \( X = -17 R_E \). For 15 August 2001, the small κ region is patchy, but it covers a much broader area in the tail, extending to about \( x = -15 R_E \), much further earthward than for 15 February 2008. The larger area with small κ implies that electrons in the 15 August 2001 event have a greater chance to experience non-adiabatic effects than those in the 15 February 2008 event.

For 15 February 2008, there was a very strong flow (~300 km/s) channel near the launch location, which represents strong \( E \times B \) drift in this region. This strong flow channel brought the launched electrons quickly from the small κ (non-adiabatic) region to the large κ (adiabatic) region. For 15 August 2001, the flows were very weak (~100 km/s) near the launch location, which means the electrons spent more time in the small κ region, and possibly near the neutral line. This also implies that electrons in the 15 August 2001 event have a greater opportunity to undergo non-adiabatic motion.

It is worth noting that at the locations with the same magnetic field radius of curvature, the higher energy electrons have smaller κ and are more likely to experience non-adiabatic effects. Figure 3-8 only shows κ for the electrons with 2-keV thermal energy. For higher energy electrons, the blue regions must be even broader. Moreover, the distribution functions result from an accumulation of non-adiabatic effects, including electric field acceleration near the weakly magnetized neutral line and scattering by the small structures in the thin plasma sheet. This can explain the differences in the high energy tail of the distribution functions between the
two events as seen in Figure 3-7. The point that electrons in the 15 August 2001 event statistically experience more non-adiabatic effects than those in the 15 February 2008 event is also confirmed by single particle calculations (not shown in this manuscript). The adiabaticity parameter distribution (Figure 3-8) demonstrates this difference in a statistical view.

3.5. Discussion

Using global MHD and LSK simulations, we have investigated the spatial evolution of electron distribution functions in the near-Earth plasma sheet for two substorm events on 15 February 2008 and on 15 August 2001.

In this study, the particles we launched were essentially near the flow channels and they generally propagated by following the flows. We did not examine the particles coming through the bursty bulk flows (BBF, discussed in Chapter 1) flanks of the flow channel as discussed by Birn et al. (2013). In order to investigate this effect, many more particles need to be launched over broader regions and that would have required massive computational resources that are beyond the scope of this study. These other source regions, however, might be important and will be considered in future studies.

As shown in the MHD simulations, the two events exhibited very different magnetotail configurations: for 15 February 2008, localized neutral lines and a narrow (< 5 \( R_E \) in \( y \) direction) earthward flow channel with relatively high speed (\( \sim 300 \) km/s) were seen in the magnetotail, while for 15 August 2001, cross-tail neutral lines and broad (> 15 \( R_E \) in \( y \) direction) earthward flows with relatively slow speed (\( \sim 100 \) km/s) were obtained. The differences were most likely caused by the dayside reconnection locations driven by solar wind conditions.

As shown in Figure 3-3, for 15 August 2001, IMF \( B_y \) component is constantly around 4 nT and dominant in the entire interval of interest. The IMF \( B_z \) component is positive (< 2 nT)
between 0800 and 0811 UT, around zero between 0811 and 0821 UT, and became negative from 0821 UT (> −2 nT). For 15 February 2008, before 0313 UT, the IMF $B_x$ component is dominant, accompanying a weak negative $B_z$ component (> −2 nT), and after 0313 UT, the IMF $B_y$ component became dominant until 0354 UT. For both events, during the intervals of interest, IMF $B_y$ components are dominant and $B_z$ components are weak.

Figures 9a and 9b show $x$-$y$, $y$-$z$ and $x$-$z$ projections of field lines at 0802:20 and 0802:40 UT for the 15 August 2001 event. The IMF, open and closed field lines are shown in red, yellow and blue respectively. The dipole tilt for this event is 7.75° toward the Sun and 27.69° toward dawn. Comparing Figure 3-9a and 3-9b, we can see that the field line topology changes in the solid rectangle and dashed rectangle regions. In the solid rectangle of the $y$-$z$ projection, there is a closed field line (blue) at 080220 UT becoming an open field line (yellow) at 080240 UT; in the dashed rectangle of the $y$-$z$ projection, there is an IMF field line (red) at 080220 UT becoming an open field line (yellow) at 080240 UT. These changes of the field line configurations imply that the reconnection happens at the locations near (6, 10, 5) $R_E$ and (10, −1, −6) $R_E$ in GSM coordinates surrounded by the solid and dashed rectangles. These reconnection regions happen at the same time and are at the high latitudes in both of northern and southern hemispheres, far away from the magnetic equatorial plane and occur because the Earth's dipole field has a negative $B_y$ component at high latitudes, which is opposite to the dominant IMF $B_y$. After the magnetic fluxes enter the magnetosphere, the dominant IMF $B_y$ component penetrates into the lobe and the plasma sheet and generates torque, tilting the plasma sheet. In the simulation, the plasma sheet is tilted about 13° (relative to the $z$ axis in GSM) to the dawn at ~20 $R_E$ in the tail. This dominant IMF $B_y$ effect of tilting the plasma sheet has been discussed by Cowley [1981]. The reconnection locations at the flanks and rotating plasma sheet
are consistent with the analysis based on the observations of this event by McPherron et al. [2002]. The penetrated IMF $B_y$ in the lobe and the plasma sheet is stretched in the $y$ direction by the solar wind and thus generates pressure in the $z$ direction to thin the plasma sheet. After the plasma sheet becomes very thin in a large region across the near-Earth magnetotail, magnetotail reconnection then can be more easily triggered in regions across a broad width of the tail. This explains the appearance of cross-tail neutral lines and patchy reconnection in a broad region in the MHD simulation for this event. In addition, the patchy reconnection in a broad region releases less energy at each site and slower bursty earthward flows than localized strong bursty reconnection. Hwang et al. [2011] also pointed out from their observations that for this event, patchy reconnection could be a reason for the observations of multiple dipolarizations, which is consistent with our simulations here. Moreover, the tilt and thinness of the plasma sheet cause very complicated and multiple scale magnetic field configurations. The complicated and multiple scales of plasma sheet during this event are also seen in the observations. For example, Weygand et al. [2005] used observations of this event as one case to discuss plasma sheet turbulence. The complicated and multiple scale magnetic configurations also explain the broad non-adiabatic regions for energetic electrons in the magnetotail during this event, as shown in Figure 3-8b.
(a) Aug 15 2001, 080220 UT

Dayside reconnection at (6,10,5) and (10,-1,-6) R_E GSM

(b) Aug 15 2001, 080240 UT

Dayside reconnection at (6,10,5) and (10,-1,-6) R_E GSM
Figure 3-9 Dayside magnetic field line configurations in $x$-$y$, $y$-$z$, $x$-$z$ projections. (a) and (b) show the field line configurations at 0802:20 and 0802:40 UT for Aug 15, 2001. (c) shows the field line configurations at 0321:00 UT for Feb 15, 2008. The IMF field lines, the open field lines and the closed field lines are shown in red, yellow and blue respectively. The arrows on the fields indicate the magnetic field direction. The solid and dashed rectangles indicate the locations of dayside reconnection. It is worth noting that in Figure 3-9c, the draped IMF field line actually provides a southward component at the nose of the magnetopause because of the large tailward dipole tilt (22.94°). The result is sub-solar reconnection for the 15 February 2008 event. In contrast, for the 15 August 2001 event, the dominant IMF $B_y$ and small sunward tilt (7.75°) yield a thin plasma sheet across the tail accompanied by reconnection over a large region in Y. See text for details.
For the 15 February 2008 event, the \( x-y \), \( y-z \) and \( x-z \) field line projections at 0321:00 UT are shown in Figure 3-9c. In the \( y-z \) projection, there are two open field lines at around (8, -3, 3) \( R_E \) GSM forming a localized reconnection configuration in the dashed rectangle, near the nose of magnetopause. Compared to the 15 August 2001 event, the dayside reconnection for this event is near the magnetic equatorial plane. Although the two events show similar IMF conditions, i.e., dominant \( B_y \) and weak \( B_z \), the dipolar tilt may be important for the difference in the reconnection locations. The dipole tilt for 15 February 2008 is 22.94° toward the tail and 22.71° toward the dusk. The IMF field lines drape around the magnetopause, which results in a sunward \( B_x \) component at the nose of the magnetopause, as shown in the square of \( x-y \) projection in Figure 3-9a and 3-9c. Because there is an apparent dipole tilt toward the tail, the sunward IMF \( B_x \) component can contribute to the reconnection near the nose of the magnetopause. This is similar to the effect that had been discussed by Russell and McPherron [1973a], in which they showed that the IMF toward (away from) the Sun had a southward component in solar magnetospheric coordinates in the spring (fall) on average. The open field lines generated by localized reconnection near the nose of magnetopause propagate to the lobe. They contribute to the relatively localized accumulation of magnetic flux in the magnetotail and further localized neutral lines and strong and narrow (in the \( y \) direction) bursty bulk flows in the magnetotail. It is worth noting that unlike the 15 August 2001 event, which had dominant IMF \( B_y \) component for a long time even before the interval of interest, for this event, the IMF \( B_y \) component became dominant starting from 0313 UT. Therefore, the plasma sheet was not twisted or as thin as seen in the 15 August 2001 event.

The complicated magnetic configuration in the magnetotail for 15 August 2001 also plays a significant role in changing the distribution function of energetic electrons. On 15 August 2001,
the distribution functions are anisotropic with \( f(v_\parallel) > f(v_\perp) \). This could be caused by adiabatic Fermi acceleration and/or non-adiabatic effects. According to Pan et al. [2012], the differential particle flux after Fermi acceleration can be estimated by multiplying a factor to the initial differential particle flux, which is Maxwellian in this study. We use the same method to calculate Fermi acceleration and find that the differential particle flux of the LSK electrons in this study has smaller slope than the estimate of Fermi acceleration. This means that Fermi acceleration cannot totally explain the field-aligned anisotropy for the 15 August 2001 event and non-adiabatic effects probably play an important role. The energetic electrons in the regions with sharply curved magnetic field lines, where the guiding center approximation fails, can experience non-adiabatic effects. In the simulation, these effects mainly include acceleration by the reconnection parallel electric field near the X-line and scattering by the small-scale electromagnetic structures in the plasmas sheet. For the 15 August 2001 event, the parallel electric field appears in a broad region accompanying the cross-tail neutral line, while for the 15 February 2008 event, the parallel electric field is only found near the localized reconnection. This means that electrons for the 15 August 2001 have more possibility to experience the parallel electric field acceleration, which causes field-aligned anisotropy in the electron velocity distribution. Parallel electric field acceleration has also been found in both observations [e.g., Øieroset et al., 2002] and other simulations [e.g., Drake et al., 2003]. The non-adiabatic scattering by the small-scale electromagnetic structures in the plasma sheet could also occur when the gyroradius of the electron is comparable to the radius of curvature of field lines, i.e., the guiding center approximation fails. For the electrons with the same energy, the ones with large pitch angle, i.e., large perpendicular velocity and gyroradius, experience non-adiabatic scattering while smaller pitch angle particles experience non-adiabatic scattering only when
they are near the equator. Most of the time, they are away from the equator. The distribution of large pitch angle electrons thus becomes more isotropic due to the scattering. The electrons with small pitch angle are not influenced by this effect as much. As a result, the distribution function shows field-aligned anisotropy due to the reduced distribution of large pitch angle electrons. It is worth noting that the higher energy electrons coincide with larger non-adiabatic regions in the magnetotail. The accumulation of such non-adiabatic effects most likely causes the field-aligned distributions shown in Figure 3-7 for the 15 August 2001 event.

Moreover, in this chapter, a new relation between non-adiabatic effects of injected electrons and BBFs has been presented. As shown in Figure 3-8, although for both the 15 February 2008 event and the 15 August 2001 event, the entries for the injected electrons are most likely at the region near the neutral lines, which is also a region where electrons can easily experience non-adiabatic effects, the BBFs associated with reconnection during the 15 February 2008 event play a role in quickly bringing these electrons earthward to the adiabatic region and reduce the amount of non-adiabatic effects experienced by these electrons. As a result, most of electrons for 15 February 2008 mainly undergo adiabatic betatron acceleration, which has also been reported based on observations [Ashour-Abdalla et al., 2011], and display the distribution functions with $f(v_\perp) > f(v_\parallel)$. In contrast, the slow earthward flows during the 15 August 2001 event provide more chance of the electrons experiencing non-adiabatic effects to change their velocity distributions.

Therefore the differences in the observations during these two substorms can be summarized as follows:
1) During the 15 August 2001 substorm, an event with neutral lines across much of the magnetotail, electrons undergo more non-adiabatic acceleration than those in the 15 February 2008 event, an event with only localized neutral lines.

2) The regions with sharply curved magnetic field lines in the plasma sheet are seen much more often in the 15 August 2001 event than in the 15 February 2008 event. Electrons in these regions are more likely to experience non-adiabatic effects.

3) The fast earthward convection flow speeds in the plasma sheet in the 15 February 2008 event quickly bring electrons earthward from the region near the reconnection site and reduce the amount of non-adiabatic acceleration experienced by these electrons.

4) The different magnetotail configurations can be caused by dayside reconnection. For 15 February 2008, the reconnection occurs at the nose of the magnetopause mainly due to a sunward IMF $B_x$ accompanying a significant tailward dipole tilt. For 15 August 2001, the reconnection occurs at the magnetospheric flanks due to a dominant IMF $B_y$.

3.6. Summary

In this chapter electron distribution functions in the near-Earth plasma sheet during two substorm events on 15 February 2008 and on 15 August 2001 are studied by using global MHD and LSK simulation. Both the MHD simulation and LSK simulations for the two events are validated through comparisons with observations by THEMIS and Cluster spacecraft. Based on the MHD simulations, the two events have very different magnetotail configurations, which are driven mainly by the different solar wind conditions and the Earth's dipole tilt for each event. For 15 February 2008, localized neutral lines and a narrow ($< 5 \text{ RE}$ in the $y$ direction) earthward flow channel with relatively high speed ($\sim 300 \text{ km/s}$) flows are observed in the magnetotail, while for 15 August 2001, cross-tail neutral lines and broad ($> 15 \text{ RE}$ in the $y$ direction)
earthward flows with relatively slow speed (~100 km/s) are obtained. Distribution functions based on an LSK simulation show that for 15 February 2008, $f(v_{\parallel}) > f(v_{\perp})$, while for the 15 August 2001 event, the distribution functions show mainly $f(v_{\parallel}) > f(v_{\perp})$. The big differences are generally caused by the variations of the non-adiabatic effects due to the different magnetotail configurations for the two events.

In addition, it is found that non-adiabatic effects need to be considered for electrons near neutral lines and in a very thin plasma sheet during a substorm. On one hand, neutral lines are highly dynamic in the magnetotail and can exist in a very broad region. On the other hand, non-adiabatic effects can accumulate to modify the distribution functions. The resulting anisotropy in the distribution functions can be unstable and generate plasma waves, which can further interact with electrons and complicate the acceleration process. Plasma wave generation and wave-particle interactions are beyond the scope of MHD and LSK simulations, however, particle-in-cell (PIC) simulations could be used to investigate these effects in future research.

In conclusion, magnetotail configurations, including flows, neutral line configuration, and plasma sheet structures, are important in determining the effect of acceleration processes in the magnetotail, particularly in the regions near neutral lines. Studying the effects of particle acceleration in the context of magnetotail configurations during different substorms is of great importance and further exploration of different events is needed in order to generalize the results found here. It would be interesting to determine the energy dependence on the non-adiabatic effects, how the energy gain is spatially and temporally related to the non-adiabatic region. It is also important to statistically study this effect and quantitatively compare the efficacy of non-adiabatic acceleration of electrons with other acceleration mechanisms. We will conduct more studies into these issues in the future to improve the understanding of this effect.
CHAPTER 4

Effects of Oxygen Ions on Magnetotail Reconnection

4.1. Introduction

Spacecraft observations show that the concentration of oxygen ions (O⁺) varies greatly between storm-time and non-storm substorms near the X-point in the magnetotail [e.g., Kistler et al., 2005, 2006]. Near the X-line, the O⁺ ions are a minor species during non-storm substorms, but they can become a major species during some of the storm-time substorms [e.g., October 1, 2001 event in Kistler et al., 2005]. The large variation in the oxygen concentration near the magnetotail X-line leads to an important question as to how much it can influence the reconnection onset, the reconnection rate and the subsequent energy transfer with propagating dipolarization fronts (DFs). There have been a number of models for the origin of the DFs such as transient structures generated by magnetic reconnection in the magnetotail [e.g., Sitnov et al., 2009; Runov et al., 2012], interchange-generated flow heads [e.g., Pritchett and Coroniti, 2011], or a $B_z$ hump motion without involving significant magnetic topology change [e.g., Sitnov et al., 2013]. In this dissertation except for Chapter 3, we only consider the DFs generated by magnetotail reconnection. In this interpretation they are the leading edge of magnetic flux in a reconnection jet and their properties are intertwined with reconnection dynamics. Using Particle-In-Cell (PIC) simulations, Wu and Shay [2012] show that DF propagation can be scaled by the Alfvén speed at the upstream edge of the ion diffusion region much as was used to scale the reconnection rate [Shay et al., 2004]. This is comparable to the plasma outflow speed [Wu et al, 2011].
In this chapter, we investigate the O$^+$ effects on the reconnection rate and the reconnection generated DFs. We used a 2.5D implicit PIC simulation of anti-parallel reconnection in the presence of H$^+$ and O$^+$ ions. The simulation mimicked magnetotail reconnection with equal number densities of O$^+$ and H$^+$. The O$^+$ ions were involved in the initial Harris current sheet [Harris, 1962] and the lobe during the reconnection. A simulation with only protons and electrons was used for comparison. The O$^+$ effects on DF propagation and reconnection rate are discussed.

4.2. Simulation Set-up

To study O$^+$ behavior without losing electron kinetic effects, an implicit PIC code, iPIC3D [Markidis et al., 2010], was used to run a 2.5D magnetic reconnection simulation. In an implicit PIC code, relatively large time steps and grid spacings can be used without loss of computational stability and energy conservation [Markidis et al., 2010]. This code has been used in several previous reconnection studies [e.g., Lapenta et al., 2010, 2013, 2015; Divin et al., 2010; Goldman et al., 2014], including a study with oxygen species present during the reconnection event [Markidis et al., 2011]. The initial state was a 2D Harris current sheet equilibrium in the X-Z plane without a guide field. The magnetic field is given by:

$$B_x(z) = B_0 \tanh \left( \frac{z - z_0}{\Delta} \right)$$

and the density is given by

$$n(z) = n_0 / \cosh \left( \frac{z - z_0}{\Delta} \right) + n_b$$

where $B_0$ is the ambient magnetic field, $n_0$ is the density at the center of the initial current sheet, $n_b$ is the background density (initially uniform in the simulation box), $\Delta$ is the half thickness of
the current sheet and $z_0$ is the Z location of the center of the current sheet. In this simulation the background density was $n_b = 0.1n_0$ for each species. Both the current sheet and the background included electron, proton and oxygen species. The density ratios with respect to the electron density were $n_e : n_H : n_O = 1 : 0.5 : 0.5$ for both current sheet and background species. This ion density ratio is typical of that observed near the near-Earth magnetotail X-line during storm-time substorms (e.g., the X-line crossing event by Cluster on August 17, 2001 [Kistler et al., 2005; Karimabadi et al., 2011]). In this study we used the Geocentric Solar Magnetospheric (GSM) coordinate system, in which the X axis points from the Earth to the Sun, the Z axis is the projection of the Earth's magnetic dipole and pointing at the geographic north direction, and the Y axis completes a right hand system. In the magnetotail current sheet, the X points at the earth, the Y points at the current direction and the Z is normal to the current sheet. In order to study both the propagation of DF and the $O^+$ characteristic scales, we used a large simulation box with $L_x = 300 \, d_{H'}$ and $L_z = 50 \, d_{H'}$ (about $23.5 \, R_E \times 3.9 \, R_E$ in the near-Earth magnetotail plasma sheet, or $116.4 \, \rho_O \times 19.4 \, \rho_O$, where $\rho_O$ is the oxygen gyro-radius). The proton inertial length is given by:

$$d_{H'} = c \left( \frac{n_0 e^2}{\varepsilon_0 m_H} \right)^{-1/2},$$

where $m_H$ is the proton mass. The current sheet density $n_0$ is given by:

$$n_0 = n_e = n_H + n_O$$

We used $d_{H'}$ in order to distinguish it from the local proton inertial length $d_H$ where

$$d_H = c \left( \frac{n_H e^2}{\varepsilon_0 m_H} \right)^{-1/2}.$$
Using the reference scale $d_H$ proved helpful when we compared the lengths in simulations with a different H$^+$ concentration setup. The simulation box included $3,840 \times 640 = 2,457,600$ cells and $\sim 2.89 \times 10^9$ ions and electrons. Mass ratios of electrons, protons and oxygen ions with respect to H$^+$ are $1/128:1:16$, respectively. Based on a statistical study by Kistler et al. [2006], the H$^+$: O$^+$ temperature ratio was approximately 1:1 near the X-line during both storm-time and nonstorm-time substorms. In this simulation, we used the temperature ratio $T_e : T_H : T_O = 0.2 : 1 : 1$ for both the current sheet and background species. The initial current sheet half width $\Delta$ was $0.5 \times d_H$ and the initial electron thermal velocity was set to $\Omega_e / \omega_s = 0.346$. The length scale was normalized by $d_H$ and the time scale was normalized by

$$\Omega_{H}^{-1} = \left( \frac{eB_0}{m_H} \right)^{-1}$$

The simulation was initially perturbed by using:

$$\delta A_y = A_{y0} \cos \left[ \frac{\pi (x - L_x / 2)}{10 \Delta} \right] \cos \left[ \frac{\pi (z - L_z / 2)}{10 \Delta} \right] e^{-((x-L_x/2)^2+(z-L_z/2)^2)/\Delta^2}$$

As in the Geospace Environment Modeling (GEM) reconnection challenge [Birn et al., 2001], this perturbation started the simulation from the non-linear regime of the tearing instability. However, in order to avoid generating initial waves with system scales, it was designed to be localized and decayed exponentially toward the boundaries [Lapenta et al., 2010]. The initial X-point was at the center of the simulation box, i.e., at $X=150$ and $Z=25$. Open boundary conditions were used for the particles in the X and Z directions. Perfect conductor boundary conditions were set at the Z boundaries for the fields. The run having equal H$^+$ and O$^+$ concentrations (i.e., $n_e : n_H : n_O = 1 : 0.5 : 0.5$) as mentioned above will be referred to as the “O$^+$ Run” in this manuscript. A run with equal mass density of H$^+$ and O$^+$ ions (i.e., $n_e : n_H : n_O = 1 : 0.5 : 0.5$) as mentioned above will be referred to as the “O$^+$ Run” in this manuscript.
0.941 : 0.059; referred to as the “EMD Run”) and a run with only H\(^+\) ions and electrons (hence referred to as the “H\(^+\) Run”) were performed and used for comparison purposes. All of the other parameters, except for the density ratio, were the same as for the O\(^+\) Run.

### 4.3. Reconnection Rate

In Figure 4-1, the reconnection rate determined for the O\(^+\) Run is compared with that for the H\(^+\) Run. The reconnection rate is estimated as the time derivative of the reconnected flux, which is calculated by taking the difference between the maximum and minimum of the Y component of the vector potential (\(A_y\)) at the center of the current sheet. In order to compare the reconnection rates of different runs, they were normalized to the proton Alfvén speed at the upstream edge of the ion diffusion region.
Figure 4-1 Reconnection Rate normalized to the proton Alfvén speed at the upstream edge of the ion diffusion region.

For the H⁺ Run, the reconnection rate reached a peak value 0.14 at $\Omega_{Ht} = 11.4$. After this point, the reconnection rate dropped to 0.1, around which it oscillated. This is a typical value for collision less reconnection [e.g., Shay et al., 1999]. The oscillations in the reconnection rate during the H⁺ Run are associated with the formation of small secondary magnetic islands (on the scale of the electron inertial length) near the X-point [e.g., Wan et al., 2008; Karimabadi et al., 2007; Bulanov et al., 1978; Loureiro et al., 2008; Lapenta, 2008].

For the O⁺ Run, fast reconnection started at around $\Omega_{Ht} = 22.0$. The reconnection rate reached a maximum value of ~0.06 at this time and then decreased gradually to 0.036 at $\Omega_{Ht}$.
This result indicates that: (1) the O\(^+\) Run reaches its fast reconnection phase later than the H\(^+\) Run, and (2) the O\(^+\) Run has a much smaller reconnection rate than the H\(^+\) Run during the fast reconnection phase.

Previous PIC simulations [Markidis et al., 2011; Karimabadi et al., 2011] showed that reconnection involving O\(^+\) goes at a slightly smaller reconnection rate than in the case without O\(^+\). Markidis et al. [2011] only considered O\(^+\) population mixed with H\(^+\) as a background plasma. Karimabadi et al. [2011] studied how the lobe plasma is assimilated into the current sheet during reconnection. They considered cases with only H\(^+\) as the initial current carrier in the presence of O\(^+\) background plasma. They also considered one case with O\(^+\) replacing H\(^+\) as an initial current carrier in order to investigate the O\(^+\) effect on the energy conversion efficiency after they are assimilated into the current sheet. The starting time of the fast reconnection phase and the reconnection rate values for our H\(^+\) Run are similar to those determined in the two previous studies. However, the reconnection rates for the O\(^+\) Run in this study developed more slowly and were much smaller later in the fast reconnection phase than those for the H\(^+\) Run and those simulations without O\(^+\) as an initial current carrier in the previous studies. In this study, O\(^+\) was an initial current sheet species, which gave the current sheet greater inertia than an H\(^+\) current sheet. A similar reconnection rate result was shown in one case with an O\(^+\) initial current carrier by Karimabadi et al. [2011], but they did not further discuss how the O\(^+\) current sheet delays the fast reconnection phase and reduces the reconnection rate. It may be more difficult for the reconnected magnetic flux to push the heavier oxygen current sheet away leading to the delay in the start of the fast reconnection phase and the much smaller reconnection rate. We will discuss this issue in more detail in Section 4.6. Since the reconnection rate is much smaller during the O\(^+\) Run, the related energy release is much slower than that observed during
the H+ Run. This moderate energy release makes the neutral sheet more stable to secondary island generation. A linear Vlasov calculation by Karimabadi et al. [2011] theoretically showed a smaller tearing growth rate in the presence of an O+ current carrier and background species. In simulations up to \(2000 \Omega_{H}^{-1}\) long Karimabadi et al. [2011] found secondary island formation was slower and there were fewer secondary islands in runs with O+ ions than in runs with only H+ ions.

4.4. Dipolarization Fronts

As noted earlier, a DF forms at the leading edge of a reconnected magnetic flux pile-up region in the outflow region. It pushes the initial current sheet in front of it. It is a layer with a steep gradient in the absolute value of \(B_z\) (from zero to a peak value at the center of the current sheet toward the initially perturbed X-point).

4.4.1. Propagation of Dipolarization Fronts

We calculated the \(B_z\) values at the center of the current sheet (along the \(X\) axis) for the O+ Run and examined their time evolution. Figures 4-2a and 4-2b show snapshots of \(B_z\) in the \(X-Z\) plane at \(\Omega_{H} t =40.2\) and \(\Omega_{H} t =89.9\), respectively. Note that the peak values of \(B_z\) (both positive and negative) correlate with the DFs. The solid black lines are magnetic field lines, calculated as contours of \(A_y\). The values of \(B_z\) along the blue and green dashed lines in Figures 4-2a and 4-2b were used to produce Figure 4-2c.
(a) $B_z$ [nT] \hspace{1cm} \Omega_H t = 40.1874

(b) $B_z$ [nT] \hspace{1cm} \Omega_H t = 89.9433

(c) $Z=24.96$ (Near the Center of the Current Sheet)
Figure 4-2 (a) $B_z$ in the $X$-$Z$ plane at $\Omega_{Ht}=40.1874$ (b) $B_z$ in the $X$-$Z$ plane at $\Omega_{Ht}=89.9433$. The solid lines indicate magnetic field lines. (c) Time evolution of $B_z$ at the center of the current sheet. The positive and negative peaks indicate the dipolarization front at a certain $X$ location at a certain time. The green and blue dashed lines in (c) correspond to the times denoted in (a) and (b), respectively. The green and blue dashed lines in (a) and (b) indicate the center of the current sheet where the $B_z$ values were taken for use in panel (c).

In the O$^+$ Run, the first X-point generated by the initial perturbation is at the center of the box ($X=150$). Later, there are two reconnection X-points generated in both the left and right outflow regions. Making a careful analysis of $B_z$ line plots at the center of the current sheet (not shown here), we found that one X-point starts at about $X=50$ when $\Omega_{Ht}=50$, while another starts at about $X=280$ when $\Omega_{Ht}=60$. The two X-points in the outflow region are possibly due to the ion tearing instability, as described by Divin et al. [2007], or the perturbation of the super-Alfvénic propagating Hall quadrupole magnetic field [Shay et al., 2011; Lapenta et al., 2013]. Due to the formation and evolution of these two X-points, two secondary islands formed on both the left and right sides of the original X-point. In this work, we mainly focused on the DF from the original X-point. We investigated its propagation and properties prior to any effects due to the growth of the secondary islands. As shown in Figure 4-2c, the peaks in $B_z$, i.e., the DFs from the original X-point, began to propagate rapidly when $\Omega_{Ht}=23$, which was approximately the time the fast reconnection phase began. After accelerating during the period $\Omega_{Ht}=23$ to $\Omega_{Ht}=35$, the speed of the DFs became roughly constant before they interacted with the secondary islands at around $\Omega_{Ht}=80$. 

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Figure 4-3 The dipolarization front speed for the O⁺ and H⁺ Runs. This speed is normalized to the heavy Alfvén speed at the upstream edge of the O⁺ diffusion region for the O⁺ Run, and normalized to the proton Alfvén speed at the upstream edge of the H⁺ diffusion region for the H⁺ Run. See text for details.

Figure 4-3 shows the time evolution of the DF speeds for the O⁺ and H⁺ Runs. These results were calculated by using the time derivative of the trace of the $B_z$ peak from the original X-point at the center of the current sheet. The DF speed was normalized to the Alfvén speed upstream of the ion diffusion region. For the O⁺ Run, the Alfvén speed at the upstream edge of the O⁺ diffusion region is given by:

$$V_{A,\text{upstream}} = B_{\text{upstream}} / \sqrt{\mu_0 (n_{H,\text{upstream}} m_H + n_{O+,\text{upstream}} m_{O+})}$$

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For the H⁺ run, the Alfvén speed at the upstream edge of the proton diffusion region is given by:

\[ V_{AH,\text{upstream}} = B_{\text{upstream}} / \sqrt{\mu_0 n_{H,\text{upstream}} m_H} \]

The region defined as the ion diffusion region is characterized by the separation of the ion flow from the \( E \times B \) flow. The normalization presented above was used because previous studies [e.g., Wu et al., 2011; Wu and Shay, 2012] showed that during a 2.5D PIC simulation, the DF speed, as well as the reconnection rate and the width of the ion diffusion region, scale with the plasma parameters at the upstream edge of the ion diffusion region.

As shown in Figure 4-3, the DF speeds for both the O⁺ and H⁺ Runs have similar patterns and magnitudes if we do not consider the time shift. This means that the DF speed scales well with the upstream Alfvén speed. As seen in Figure 4-3, the DF speed increased with time for both runs. After the start of the fast reconnection phase (\( \Omega_{Ht} = 23 \) for the O⁺ Run; \( \Omega_{Ht} = 12 \) for the H⁺ Run), the DF speed increased in two phases. For the O⁺ Run, the DF acceleration was large during the period \( \Omega_{Ht} = 23 \) to \( \Omega_{Ht} = 40 \), but became much smaller after \( \Omega_{Ht} = 40 \). For the H⁺ Run, the rapid increase phase occurred from \( \Omega_{Ht} = 12 \) to \( \Omega_{Ht} = 20 \). Later, the acceleration decreased but was still much greater than that in the same phase for the O⁺ Run. During the second phase, the DF speed in the O⁺ Run increased during the period 0.5 \( V_{A,\text{upstream}} \) to 0.6 \( V_{A,\text{upstream}} \), while it increased during the period 0.5 \( V_{AH,\text{upstream}} \) to 0.7 \( V_{AH,\text{upstream}} \) in the H⁺ Run.

For a two-species 2D anti-parallel reconnection Wu and Shay [2012] measured the DF speeds in simulations using different background densities. They found that the DF speeds scaled well with the upstream Alfvén speed, with a linear fit given by \( V_{DF} \approx 0.3 V_{AH,\text{upstream}} \). It is worth noting that the parameters used by Wu and Shay [2012] are different from those in our H⁺ Run. They used the mass ratio \( m_i/m_e = 25 \), set the half thickness of the initial current sheet to be \( \sim 1 \)
and used periodic boundary conditions. We used $m_H/m_e = 128$, set the half thickness to 0.5 $d_H$, and used open boundary conditions.

Compared with the H$^+$ Run, the DF speed in the O$^+$ Run takes about twice as long to evolve from the first phase through the second phase. Note that the DF speed in the O$^+$ Run is normalized to an upstream Alfvén speed with O$^+$ mass so it would be much smaller than that in the H$^+$ Run if they both were normalized to the local proton Alfvén speed. The slow propagation of the DF in the O$^+$ Run was probably due to the large inertia of the initial current sheet, which contained an O$^+$ species. We will examine this point in the following sections by analyzing the force balance and plasma composition at the DF position.

4.4.2. Force Balance at a Dipolarization Front

In order to understand the propagation of a DF, we analyzed the force balance at its position. The DFs analyzed in both the O$^+$ and H$^+$ Runs occurred during the second phase. The forces determined at each DF position included the thermal pressure gradient force due to each plasma species and the stress force due to the electromagnetic field. The elements of the thermal pressure tensor for each species were calculated as the second moment of the velocity distribution function:

$$P_{ij}^{(s)} = \int m^{(s)}(v_i^{(s)} - <v_i^{(s)}>)(v_j^{(s)} - <v_j^{(s)}>)f^{(s)}(\tilde{v})d^3\tilde{v}$$

where the superscript $s$ indicates different species and the subscripts $i$ and $j$ indicate which $x$, $y$, or $z$ components were used. The thermal pressure gradient force density was determined as the negative divergence of the pressure tensor $-\nabla \cdot \vec{P}^{(s)}$. The stress force density of the electromagnetic field is calculated by using $\nabla \cdot \vec{T}$, where $\vec{T}$ is the Maxwell tensor:
\[ T_{ij} = \varepsilon_0 E_i E_j + \frac{1}{\mu_0} B_i B_j - \frac{1}{2} \left( \varepsilon_0 E^2 + \frac{1}{\mu_0} B^2 \right) \delta_{ij} \]

Figure 4-4 (panels a-c) shows the \( B_z \) component, the \( X \) component of the thermal pressure gradient force density of each species, the force balance between the electromagnetic stress forces, and the total thermal pressure gradient force in the \( X \) direction for the \( O^+ \) Run when \( \Omega_{Hi} = 61.238 \). All of these quantities were measured 0.04 \( d_{Hi} \) (minimum half grid size) from the center of the current sheet along the \( X \)-axis. Note that the X-point is at \( X = 150 \) and the DF shown in Figure 4-4a is propagating to the right, in the positive \( X \) direction. The positive (negative) force density corresponds to the positive (negative) \( X \) direction. The DF is the layer from the \( B_z \) peak value to zero as delineated by the two vertical dashed lines in Figure 4-4 (panels a-c). As shown in Figure 4-4b, the \( O^+ \) species was the major contributor to the force balance in the DF region, while the \( H^+ \) species barely made a contribution. However, electrons created an opposite force in that region, which was due to their distribution function. We will discuss this issue later in Section 4.6. Figure 4-4c shows the transient force balance between the electromagnetic stress force and the total thermal pressure gradient force. The strongest force density occurs at the center of the DF.
Figure 4-4 (a)-(c): $B_z$, pressure gradient force density for electrons, H$^+$ and O$^+$ ions, Maxwell stress tensor force density and total pressure gradient force density for the O$^+$ Run. (d)-(f): $B_z$, pressure gradient force density for electrons and H$^+$, Maxwell stress tensor force density and total pressure gradient force density for the H$^+$ Run. All these quantities are measured at the center of the current sheet along the X-axis. The DF layer is highlighted by the double vertical dotted lines for both runs. See text for details.

For comparison, the force balance at the DF in the H$^+$ Run when $\Omega_{Ht}=30.619$ is shown in Figure 4-4 (panels d-f) in the same format as for the O$^+$ Run. At that moment, the DF was the same distance from the center X-point as it was in the O$^+$ Run. As seen in Figures 4-4e and 4-4f, H$^+$ was the major contributor to the force balance, while the electrons in that run did not exert any significant force at the DF region. Note that the ions may not be frozen-in near the front part of the DF layer for both runs due to the weak magnetic field. As we will show in Section 4.4.3, the thickness of the DF layer was similar to the local O$^+$ inertial length (the local H$^+$ inertial length) for the O$^+$ Run (the H$^+$ Run).

As will be shown in Section 4.4.3, the ion force on the DF is caused by the initial current sheet ion species. The heavier ion mass of the current sheet in the O$^+$ Run leads to a smaller acceleration of the DF compared to the H$^+$ Run. It should be noted that the force balance that occurred during the first phase of the DF propagation was similar to that occurring in the second phase.

4.4.3. Thickness and Plasma Composition of the Dipolarization Front

To understand how the heaviest ion species produces the major force in the DF region, we examined the thickness of the DF and its plasma composition. Figure 4-5 (panels a-c) shows
the $B_z$ component, the local ion gyro-radius $\rho$, the ion inertial length $d$ for O$^+$ and H$^+$, and the density for each species including the initial current sheet (“0”) and background (“b”) components for the O$^+$ Run when $\Omega_{Ht} = 61.238$. All of these quantities were measured 0.04 $d_H$ (minimum half grid size) away to the center of the current sheet along the $X$-axis. The current sheet component included the drifting plasma set up as the initial current sheet, which was expelled by the DF. Initially, the background component covered the entire simulation box (including the current sheet region) and was mainly comprised of the lobe plasma species. Figure 4-5a shows that the DF had a thickness of 2.6 $d_H$. The dotted lines in Figure 4-5b show the extent of the DF. The DF thickness was close to the local O$^+$ inertial length $d_O$ in the DF region and also comparable to the local $d_H$ even though $d_H$ varied significantly in this region. Note that $\rho_H$ is small in front of the DF even though $B_z$ is close to zero. This is because the point of reference is not exactly at the center of the current sheet and thus both $B_x$ and $B_y$ are not close to zero. The $B_y$ in front of the DF is generated by reflected ion current [e.g., Wu and Shay, 2012]. In Figure 4-5c, the current sheet ions (O$^+$ and H$^+$) show different penetration depths in the DF. Because of the greater O$^+$ penetration (mainly due to the larger gyroradius at the front edge), the number density component of the current sheet O$^+$ (red solid line) was dominant at the DF and this mainly determined the thickness of the DF. The current sheet's electron component partly compensated for the penetrating O$^+$ and H$^+$ ions. A significant background electron component compensated for the overshooting O$^+$ ions near the location with the peak $B_z$ value. The current sheet O$^+$ also provided the major force balance at the DF as shown in Figure 4b and 4c. Note that the simulation started with a very thin current sheet that mimicked a thin plasma sheet before the onset of magnetotail reconnection. Also note that such a thin long (300 $d_H$) uniform current sheet is unlikely in the real magnetotail due to the development of tearing-like
instabilities. We used this ideal current sheet in order to examine the DF propagation in the presence of O\(^+\), while avoiding affects due to the simulation box boundaries.
Figure 4-5 (a)-(c): $B_z$, local inertial length $d$ and local gyro-radius $\rho$ of H$^+$ and O$^+$ ions, current sheet density (“0”) and background density (“b”) for electrons, H$^+$ and O$^+$ ions for the O$^+$ Run. (d)-(f): $B_z$, local inertial length $d$ and local gyro-radius $\rho$ of H$^+$, current sheet density (“0”) and background density (“b”) for electrons and H$^+$ ions for the H$^+$ Run. All these quantities are measured at the center of the current sheet along the $X$-axis. The DF layer is highlighted by the double vertical dashed lines with its thickness on top for both runs. The thickness of the DF is also shown by dashed-and-dotted lines in (b) for the O$^+$ Run and (e) for the H$^+$ Run.

For comparison, the thickness and plasma composition of the DF region in the H$^+$ Run when $\Omega_{Ht}$=30.619 are shown in Figure 4-5 (panels d-f). The thickness of the DF was about 1.2 $d_{H'}$, which is comparable to the local H$^+$ gyro-radius $\rho_{H}$ and the local H$^+$ inertial length $d_{H}$ as shown in Figure 4-5e. In Figure 4-5f, the current sheet H$^+$ density was dominant in the DF region and it determined the thickness of the DF. Thus, the force balance at the DF shown in Figures 4-4e and 4-4f was mainly provided by the penetrating current sheet H$^+$ ions. We also found that electrons provided different force contributions in the DF regions for the O$^+$ and H$^+$ Runs. We will discuss this further in Section 4.6.

In general, our simulation determined that the force at the DF is mainly provided by the current sheet's heaviest ion species. This means that DF propagation must overcome the inertia of the initial current sheet, which explains the slower propagation of the DF in the O$^+$ Run. In addition, since the outflow flux must overcome the current sheet inertia, this may explain the delay of the fast reconnection phase for the O$^+$ Run (see Figure 4-1).
4.5. Ambipolar Electric Field

An ambipolar electric field was found in the O+ Run as a result of the O+ effect on reconnection. Figure 4-6 (panels a-c) shows the $B_z$ component, the plasma outflow velocities, and the $E_x$ component at the center of the current sheet along the $X$-axis for the O+ Run when $\Omega_{Ht} =61.238$. Figure 4-6a shows the DF location at that moment. Note that the X-point was at $X=150$ and the DF was propagating to the right. Figure 4-6b compares the outflow velocities of the electron, H+, and O+ species, the $X$ component of the $E\times B$ velocity, and the DF speed. The velocities are normalized to the upstream Alfvén speed. They have identical values at the DF location (i.e., the peak value of $B_z$). However, in the outflow region behind the DF ($X<172$), the O+ outflow velocity was much less than those of the electrons and H+ ions because that region was within the O+ diffusion region. The difference between the O+ and electron outflow speeds produced an ambipolar electric field along the positive $X$ direction, i.e., the positive $E_x$ component shown in Figure 4-6c. For comparison, Figure 4-6 (panels f-h) shows the same variables (except for the O+ outflow) for the H+ Run when $\Omega_{Ht} =37.316$. As shown in Figure 4-6g, the plasma outflow velocity and the $E\times B$ velocity were the same as the DF speed during that moment. Even though the H+ outflow speeds were slightly less than the electron outflow speed in a part of the H+ diffusion region ($X<176$), we determined that in general, the H+ and electron outflow speeds were similar in the outflow region behind the DF ($X<185$). No ambipolar electric field, $E_x$, was noticed in this region for the H+ Run (see Figure 4-6h). As shown in Figures 4-6c and 4-6h, there are bipolar signatures in $E_x$ at the DFs for both the O+ Run and the H+ Run. To explain this signature, we calculated the generalized Ohm’s law for both runs. For the H+ Run, the 2-species generalized Ohm’s law is:
\[
\tilde{E} = \left( \frac{n_H}{n_e} \right) \tilde{v}_H \times \tilde{B} + \frac{\tilde{J} \times \tilde{B}}{n_e e} - \frac{\nabla \cdot \tilde{P}_e}{n_e e} - \frac{m_e}{e} \left( \frac{\partial}{\partial t} \tilde{v}_e + \tilde{v}_e \cdot \nabla \tilde{v}_e \right),
\]

where \( \tilde{J} = e n_H \tilde{v}_H - en_e \tilde{v}_e \) and \( \tilde{P}_e \) is the electron pressure tensor. It includes an H\(^+\) convection term, a Hall current term, an electron pressure tensor term and an electron inertia term. Similarly, the generalized Ohm’s law for the O\(^+\) Run is

\[
\tilde{E} = \left( \frac{n_H}{n_e} \right) \tilde{v}_H \times \tilde{B} - \left( \frac{n_O}{n_e} \right) \tilde{v}_O \times \tilde{B} + \frac{\tilde{J} \times \tilde{B}}{n_e e} - \frac{\nabla \cdot \tilde{P}_e}{n_e e} - \frac{m_e}{e} \left( \frac{\partial}{\partial t} \tilde{v}_e + \tilde{v}_e \cdot \nabla \tilde{v}_e \right),
\]

where \( \tilde{J} = e n_H \tilde{v}_H + en_O \tilde{v}_O - en_e \tilde{v}_e \). Compared with the 2-species case, it has an additional O\(^+\) convection term. Figures 4-6d and 4-6i show the X-component of all these terms in the generalized Ohm’s law compared with \( E_x \) for the O\(^+\) Run and the H\(^+\) Run, respectively. For both runs, the Hall current term was canceled by the convection terms, and the bipolar signature of \( E_x \) at the DFs was essentially determined by the electron pressure tensor term (green line in the panels). Because the electron pressure tensor term includes \( -\nabla \cdot \tilde{P}_e \) in the 2D case, we show the contributions of \( -\partial P_{e,xx}/\partial x \) and \( -\partial P_{e,xx}/\partial z \) in Figures 4-6e and 4-6j for the O\(^+\) Run and the H\(^+\) Run, respectively. As shown in Figures 4-6e and 4-6j, the \( -\partial P_{e,xx}/\partial x \) term mainly determined the bipolar signatures even though \( -\partial P_{e,xx}/\partial z \) provided a significant contribution at the DFs. This indicates that the bipolar signature was due to a pile-up electron pressure \( P_{e,xx} \) at the DF region.
Figure 4-6 (a)-(e): $B_z$, outflow velocities for electrons, H$^+$ and O$^+$ ions with the $X$ component of the $E \times B$ velocity and the DF speed at this moment, $E_x$, for the O$^+$ Run. (d) generalized Ohm’s law, including H$^+$ convection term (blue), O$^+$ convection term (red), Hall current term (pink), electron pressure tensor term (green), electron inertia term (orange), compared with $E_x$ (black), for the O$^+$ Run. (e) $-\partial P_{e,xx} / \partial x$ and $-\partial P_{e,zz} / \partial z$, for the O$^+$ Run. (f)-(j): parameters for the H$^+$ Run in the same format as (a)-(e). All these quantities are measured at the center of the current sheet along the $X$-axis. The DF layer is highlighted by the double vertical dashed lines for both runs. The red ellipse highlights the non-zero ambipolar electric field.

For the inflow region, Figure 4-7 (panels a-c) shows the $B_x$ component, the plasma inflow velocities, and the $E_z$ component along the $Z$-axis cutting through the X-point for the O$^+$ Run.
As shown in Figure 4-7b, the separation of the inflow velocity from the $E \times B$ velocity determined the diffusion region for each species. The $O^+$ diffusion region was much larger than those of the other two species. Inside the $O^+$ diffusion region, but outside the $H^+$ diffusion region ($Z<23$ or $Z>28$), the $O^+$ inflow speed was slightly less than those of the $H^+$ and electron inflows. As a result, an ambipolar electric field, $E_z$, was created in the inflow region as shown in Figure 4-7c. Note that this inflow ambipolar electric field had a much smaller value than that in the outflow region (Figure 4-6c). This was due to the difference between the plasma flow speeds in the inflow and outflow regions. The former was much less than the latter. For comparison, Figure 4-7 (panels d-f) shows the same variables (except for the $O^+$ inflow) for the $H^+$ Run when $\Omega_{nit} = 37.316$. There were no differences between the inflow speeds outside the $H^+$ diffusion region so no ambipolar electric field formed in the inflow region. In the $H^+$ diffusion region, $23 < Z < 27$ for both runs, there was a bipolar Hall electric field due to the decoupling of $H^+$ and electrons [e.g., Hoshino et al., 2001]. The ambipolar electric field in the inflow region of the $O^+$ Run could also be understood as a Hall electric field in the $O^+$ diffusion region.
Generally, the ambipolar electric field in the O⁺ Run was due to an O⁺ effect on reconnection. It was generated by the difference between the O⁺ and electron flow speeds in the inflow and outflow regions. As we will discuss in the following section, the creation of this ambipolar electric field provided another explanation for the slow reconnection rate in the O⁺ Run.
4.6. Discussion

4.6.1. Equal Mass Density Run

As shown in the previous Sections 4.4.2 and 4.4.3, the current sheet's heaviest ion component provided the major contribution to the force balance at the DF. For the $O^+$ Run (where $n_O = n_H$), the current sheet $O^+$ ions primarily determined the thickness of the DF, which was similar to the local $O^+$ inertial length $d_O$. The DF propagation had to overcome the inertia of current sheet $O^+$ ions. If the force contribution was entirely determined by the mass density gradient at the DFs, then one could speculate that if the $O^+$ and $H^+$ ions had equal initial mass densities, then they probably contributed equally to the force balance in the DF region. To examine this issue, we set up a run with equal mass densities of $H^+$ and $O^+$ ions, i.e., the “EMD Run”. Figure 4-8 shows the reconnection rates for the EMD, $O^+$, and $H^+$ Runs. The reconnection rate curve for the EMD Run was similar to that for the $H^+$ Run, but the fast reconnection phase occurred slightly later and the reconnection rate was smaller during the quasi-steady reconnection phase. These differences are understandable because the results are similar to an $H^+$ Run with a very low concentration (~5.9%) of $O^+$ ions. Figure 4-8 (Panels b-d) shows the $B_z$ component, the plasma pressure gradient force density and the force balance near the DF region. Figures 4-8e and 4-8f show the characteristic scales and plasma composition as seen in Figures 4-5b and 4-5c. Figures 4-8c and 4-8d show that the force balance was mainly provided by the $H^+$ species at the DF. Similar to the $H^+$ Run, the thickness of the DF was comparable to the local $H^+$ gyro-radius and the local $H^+$ inertial length. Figure 4-8f shows that the penetrating current sheet $H^+$ ions dominated the ion number density at the DF. Therefore, they determined the thickness of the DF and their inertia had a major effect on the DF region.
Figure 4-8 Results from the simulation with equal mass densities of O+ and H+ ions (EMD Run). (a) Reconnection rate (b) $B_z$ (c) pressure gradient force density of electrons, H+ and O+ ions (d) Maxwell stress tensor force density and total pressure gradient force density (e) local inertial length $d$ and local gyro-radius $\rho$ of H+ and O+ ions (f) current sheet density (“0”) and background density (“b”) for electrons, H+ and O+ ions. All these quantities in (b)-(f) are measured at the center of the current sheet along the $X$-axis. The DF layer is highlighted by the double vertical dotted lines with its thickness on top. The DF thickness is also indicated in (e) as a horizontal dashed-and-dotted line.

In general, the results of the EMD Run were very similar to those of the H+ Run due to the very low concentration of the O+ species. The O+ density at the DF was too low to produce any significant force in the DF region or to determine the thickness of the DF despite it being the heaviest ion species in this run. The concentration of the heaviest ion species needed to provide a significant contribution to the force balance at the DF remains undetermined. It is the pressure gradient that determines the force contribution at DFs. This is related to the ion velocity distribution at the DF, which includes reflection, heating and acceleration of the ions. This issue will be further discussed in a future study.

4.6.2. Electron Force at the DF

As shown in Figure 4-4b (O+ Run), electrons provide a significant force in the DF layer that opposes the O+ force. In Figure 4-4e (H+ Run), the electron force has a narrow positive peak at the same location as the maximum $B_z$, and it oscillates between positive and negative values in the DF layer. What causes these differences in electron force behavior? To answer
this question, we examined the electron velocity distribution function, which was used to calculate the pressure tensor. The electron pressure gradient force was calculated as:

\[
\left(-\nabla \cdot \vec{P}\right)_x = -\frac{\partial P_{xx}}{\partial x} - \frac{\partial P_{xz}}{\partial z}
\]

in the 2.5D simulation. In the DF region, \(B_z\) was dominant so \(X\) and \(Z\) corresponded to the in-plane perpendicular and parallel directions with respect to the magnetic field. Here, we define the parallel direction to the magnetic field \(\hat{b}\) as the “||” direction and the in-plane direction perpendicular to \(\hat{b}\) as the “\(\perp\)1” direction, where:

\[
\hat{e}_{\perp 1} = \frac{\hat{e}_y \times \hat{b}}{|\hat{e}_y \times \hat{b}|}
\]

The other perpendicular direction “\(\perp 2\)” is out-of-plane and follows the right hand rule. Thus, near the peak of \(B_z\) at the DF, the \(X\) and \(Z\) directions correspond to “\(\perp 1\)” and “||” directions in the equation above. Note that the approximation of “\(x\)”~ “\(\perp 1\)” and “\(z\)”~“||” is only good in the region near \(B_z\) peak of the DF. It is not a good approximation at other regions without dominant \(B_z\). We examined the electron velocity distribution function in the DF region as a function of \(V_{\perp 1}\) and \(V_{||}\) to explain the different electron forces at the DF for the two runs. The results are shown in Figures 4-9 (O+ Run) and 4-10 (H+ Run).
Figure 4-9 (a)-(c): $B_z$, current sheet (“0”) and background (“b”) density of electrons, electron temperature in the “$||$”, “$\perp 1$” and “$\perp 2$” directions, for the O+ Run. (d)-(h): electron velocity distribution functions at different $X$ locations as indicated by the blue dashed lines in (c). The DF layer is highlighted by the double vertical dotted lines. See text for details.

Figure 4-9 (panels a-c) shows the $B_z$ component, electron density, and electron temperature as a function of $X$ near the DF region. This DF is the same as that in Figure 4-4 and is highlighted by the two dotted vertical lines. As shown in Figure 4-9b, the current sheet component $N_e(0)$ and background component $N_e(b)$ are both significant in the DF layer. As mentioned previously in Section 4.4.3, the current sheet component is the species forming the initial current sheet, and the background component initially covers the entire simulation box and basically mimics the lobe plasma. As seen in Figure 4-9b, the former mainly occupied the front region of the DF
layer, while the latter dominated the region near the peak $B_z$ value for that layer. We mentioned in Section 4.4.3 that both the background and current sheet electron components compensated for the effect of the penetrating current sheet ions. The background electrons in the DF layer mainly adjusted in response to the deep penetration of the $O^+$ ions. Figure 4-9c shows the electron temperature in the “||”, “⊥1”, and “⊥2” directions, which we used to examine the anisotropy of the electron velocity distribution function. In order to compare with spacecraft observations and other studies, the temperature was calculated by using the equation of state:

$$\vec{P} = n \vec{T}$$

where $\vec{P}$ and $n$ are the pressure tensor and density, respectively. As shown in Figures 4-9c and 4-9b, $T_{e,⊥1} \approx T_{e,⊥2}$ was true everywhere. An anisotropy in which $T_{e,⊥1} > T_{e||}$ was found in the region with a dominant background component, while an isotropic distribution ($T_{e||} \approx T_{e,⊥1} \approx T_{e,⊥2}$) was found in the region with a dominant current sheet component. These electrons could experience multiple-step acceleration in different regions near the reconnection site [e.g., Hoshino et al., 2001]. We examined the electrons behind the DF when the DF leaves the electron jet. From that point, the electrons underwent $E \times B$ drift and were trapped by the magnetic field and parallel electric potential and experienced betatron acceleration while moving with the propagating DF. A similar scenario was noted by Huang et al. [2015]. The regions that show a large anisotropy ($X<175$) in Figure 4-9c correspond to the regions in which the background electron population was much larger than the current sheet electron population, as shown in Figure 4-9b. Thus, the electron anisotropy in these regions resulted from the multiple-step acceleration of the background electron species. The current sheet electrons did not experience any effective acceleration before reaching the DF. The distribution functions in
Figure 4-9 (panels d-h) correspond to the locations in the DF layer highlighted by the blue dashed lines. These distribution functions confirm the anisotropic and isotropic signatures seen in the DF layer in Figure 4-9c.

As shown in Figure 4-4b, the electrons provided a positive force in the DF layer opposite to that of the O$^+$ force between $X=175$ and $X=176$. We also see that this force was actually caused by an anisotropic distribution function involving $T_{e,\perp} > T_{e,\parallel}$. 
Figure 4-10 (a)-(c): $B_z$, current sheet (“0”) and background (“b”) density of electrons, electron temperature in the “||”, “\ perpendicular” and “\ perpendicular2” directions, for the H+ Run. (d)-(h): electron velocity distribution functions at different $X$ locations as indicated by the blue dashed lines in (c). The DF layer is highlighted by the double vertical dashed lines. See text for details.

The types of plots in Figure 4-9 are repeated in Figure 4-10 using results from the H+ Run. In Figure 4-10 (panels b and c), the current sheet electron component dominated the DF layer and the region behind the DF was dominated by the background component corresponding to the $T_{e,\parallel} > T_{e,\perpendicular}$ anisotropy. There was no significant $T_{e,\parallel} > T_{e,\perpendicular}$ anisotropy in the DF layer, but there was a region between $X=175.6$ to $X=176.0$ where $T_{e,\parallel} < T_{e,\perpendicular}$. From the distribution function in Figure 4-10g, we see a counter-streaming signature along the magnetic field.
direction corresponding to the peak ratio \( \frac{T_{e\parallel}}{T_{e\perp}} \) in the DF layer, which resulted from a localized bipolar parallel electric field at the DF (not shown in this manuscript). This localized bipolar electric field probably formed due to the instability of the current sheet. The electron force at the DF in the O\(^+\) Run was different from that in the H\(^+\) Run due to the lack of an anisotropic distribution in the DF layer.

In summary, the background electrons in the O\(^+\) Run with \( T_{e\perp} > T_{e\parallel} \) compensated for the deep penetration of the O\(^+\) ions into the DF layer, resulting in a pressure gradient force that opposed the O\(^+\) force at the DF.

4.6.3. Reconnection Rate and Ambipolar Electric Field

As discussed previously in Section 4.3, the reconnection rate in the O\(^+\) Run (1) showed a delayed start of the fast reconnection phase, and (2) was much smaller during the fast reconnection phase than in the H\(^+\) Run. The results in Section 4.4 showed that the outflow propagating flux had to overcome the inertia of the initial current sheet. During steady reconnection, the reconnection rate equaled the rate of flux inflow, which was balanced by the rate at which the reconnected flux was expelled. Since the reconnected outflow flux was encumbered by the heavy O\(^+\) ions in the initial current sheet for the O\(^+\) Run, the reconnection rate was reasonably smaller than that for the H\(^+\) Run. This also explains the delay of the fast reconnection phase for the O\(^+\) Run.

In Section 4.5, we reported evidence that ambipolar electric fields were present in the outflow and inflow regions during the O\(^+\) Run. The presence of an ambipolar electric field could explain the small reconnection rate. As diagramed in Figure 4-11, electrons followed the motion of the magnetic field, while O\(^+\) ions were demagnetized and fell behind in the O\(^+\) diffusion.
region. The difference between the $O^+$ and electron flow speeds in that region caused the ambipolar electric field. The ambipolar electric field had to accelerate the heavy $O^+$ ions in that region. That drag was actually exerted by the $O^+$ ions on the propagating magnetic flux via the magnetized electron flow. Since the reconnection rate was equivalent to the speed of the magnetic flux propagation from the inflow region to the outflow regions, its magnitude could be deduced from the $O^+$ drag via the ambipolar electric field.

![Diagram](image)

**Figure 4-11** Relationships of the $O^+$ ions, the ambipolar electric field, and the magnetic flux propagation in the $O^+$ diffusion region.

The small reconnection rate noticed in our simulations in the presence of $O^+$ was determined also in previous studies [e.g., Shay and Swisdak, 2004; Markidis et al., 2011; Karimabadi et al., 2011]. Using a fluid simulation, Shay and Swisdak [2004] showed that the reconnection rate with $O^+$ has an asymptotic value less than that of a two-species reconnection by a factor of 1.5. Markidis et al. [2011] used an implicit PIC simulation with periodic boundary conditions in $X$ that included only background $O^+$ species (no $O^+$ ions in the initial current sheet). They showed that the $O^+$ and H$^+$ Runs had identical starting times for the fast reconnection.
phase, while the reconnection rate for an O\textsuperscript{+} Run was slightly less than that of an H\textsuperscript{+} Run. That the fast reconnection starts at the same time in both cases was due probably to the lack of any heavy ion inertia in the initial current sheet in their O\textsuperscript{+} run. The smaller reconnection rate value could also be explained by considering the ambipolar electric field mechanism discussed in this work. Karimabadi et al. [2011] used PIC simulations with open boundary conditions and a mass ratio \(m_{O}/m_{H}=9\). Their runs included O\textsuperscript{+} and H\textsuperscript{+} species. The O\textsuperscript{+} ions were only set as background species initially, but there was one run with pure O\textsuperscript{+} that included the O\textsuperscript{+} species in the initial current sheet. In their runs without O\textsuperscript{+} in the initial current sheet, the fast reconnection phase started at about the same time as in the corresponding pure H\textsuperscript{+} runs (see Figure 8 in Karimabadi et al. [2011]). Using the same normalization, the reconnection rate of their O\textsuperscript{+} run was slightly higher than that in our O\textsuperscript{+} Run. This was due to the smaller mass of O\textsuperscript{+} used in their run. For their pure O\textsuperscript{+} run that included O\textsuperscript{+} ions in the initial current sheet, the reconnection rate was significantly delayed compared to the fast reconnection phase in the pure H\textsuperscript{+} run (see Figure 9a in Karimabadi et al. [2011]). This supports our hypothesis regarding the effect of the O\textsuperscript{+} current sheet inertia on the reconnection rate.

In summary we determined two factors that contribute to the delay of the fast reconnection phase and the slow reconnection rate in the presence of O\textsuperscript{+} species:

1. The outflow reconnected flux propagation is encumbered by the heavy current sheet O\textsuperscript{+} inertia.
2. The propagation speed of the magnetic flux is reduced by the drag of the O\textsuperscript{+} ions via the ambipolar electric field.
4.7. Conclusions

In this study we used a 2.5D implicit Particle-in-Cell simulation (iPIC3D) with open boundary conditions in the presence of H\(^+\) and O\(^+\) ions to investigate the O\(^+\) effects on dipolarization fronts and reconnection rate during magnetotail reconnection. An O\(^+\) Run with equal number densities of O\(^+\) and H\(^+\) and an H\(^+\) Run with only H\(^+\) and electrons were analyzed in detail. Initially, the current sheet was set up as a Harris current sheet, and the O\(^+\) species were included in both the initial current sheet and the background for the O\(^+\) Run. Using the same normalization, the DF speed in the O\(^+\) Run was much less than that in the H\(^+\) Run. By examining the force balance and plasma composition in the DF region, we found that DF propagation was mainly encumbered by the heaviest ion in the current sheet in our simulations for a sufficiently large concentration of O\(^+\). This effect caused a delay in the fast reconnection phase and a low reconnection rate. In addition, the heaviest ions in the current sheet penetrated the DF and determined its thickness. A difference between the O\(^+\) and electron flow speeds generated an ambipolar electric field in the O\(^+\) diffusion region. This ambipolar electric field caused O\(^+\) drag on the propagating magnetic flux and thus reduced the reconnection rate. Our main conclusions are:

1. The DF propagation has to overcome large current sheet O\(^+\) inertia if the current sheet includes a sufficient concentration of O\(^+\) species.

2. If the concentration of the heaviest current sheet ions is sufficient, the thickness of a DF is mainly related to the penetration of those ions. The dipolarization front thickness was comparable to the local H\(^+\) inertial length \(d_H\) and the local H\(^+\) gyro-radius \(\rho_H\) in the H\(^+\) Run. In the O\(^+\) Run, it was close to the local O\(^+\) inertial length \(d_O\) and also
comparable to the $d_H$ in the O$^+$ Run, even though $d_H$ significantly varied in the DF region.

(3) There are two phenomena that contribute to the delay of the fast reconnection phase and the slow reconnection rate in our simulations in the presence of O$^+$ ions:

a) Outflow flux propagation was encumbered by the current sheet O$^+$ inertia.

b) Magnetic flux propagation was impeded by O$^+$ ions due to the formation of an ambipolar electric field in the O$^+$ diffusion region.

We conclude that the heaviest current sheet ions provide a major force in the DF region and determine the DF thickness. However, a limitation still exists. In the previous Section 4.6.1, we studied a simulation having equal mass density of O$^+$ and H$^+$ species (i.e., $n_e : n_{H} : n_{O} = 1 : 0.941 : 0.059$). In that case, we found that the relatively low density of O$^+$ provided only a small force contribution at the DF. Therefore, the concentration of the heaviest current sheet ion is a determining factor regarding force balance at the DF. However, the concentration of a particular current sheet ion sufficient to cause a significant force at the DF is still uncertain. Further analysis regarding this issue will be in a future study.

Finally, we found that there was an ambipolar electric field in the O$^+$ diffusion region caused by the difference in the O$^+$ and electron flow speeds. This ambipolar electric field possibly caused O$^+$ drag on the propagating magnetic flux and thus reduced the reconnection rate. This effect still needs to be determined quantitatively in order to find out its actual contribution to the decrease in the reconnection rate. In addition, as observed by spacecraft, the concentration of oxygen ions (O$^+$) varies greatly between storm-time and non-storm substorms near the magnetotail X-line. The concentration of O$^+$ in the O$^+$ Run and the EMD Run could represent storm-time and non-storm substorms, respectively. The delay of the starting time of
the fast reconnection phase and the slower reconnection rate and propagation of DFs were observed in the O\(^+\) Run compared with the EMD Run and the H\(^+\) Run. This illustrates the difference between the two types of substorms. Other issues remain, including the influence of the O\(^+\) concentration on this effect, the role of the H\(^+\) ions on this effect, and the determination of any similar effect due to H\(^+\) ions in the H\(^+\) diffusion region. More detailed investigations are planned for future research.
CHAPTER 5

Ion Acceleration in magnetotail reconnection in the presence of oxygen ions

5.1. Introduction

Motivated by the observed high concentration of oxygen ions in the magnetotail during enhanced geomagnetic activity, the ion acceleration (mainly oxygen acceleration) in magnetotail reconnection is investigated in this chapter.

Since the oxygen ions can affect the reconnection process, oxygen acceleration and heating near the reconnection site must be considered self-consistently. Because of their large gyroradii, the region where oxygen ions are demagnetized is larger than that for protons. Using a PIC simulation, Liu et al. [2015] showed that in the oxygen diffusion region, the oxygen ions are heated in the exhaust and their velocity distribution functions show a counter-streaming signature due to meandering motions across the current sheet combined with a drifting beam along the $Y_{GSM}$ direction due to acceleration by the reconnection electric field. In the magnetotail current sheet, the $X_{GSM}$ points at the earth, the $Y_{GSM}$ points at the current direction and the $Z_{GSM}$ is normal to the current sheet. According to Liu et al. [2015], the features of oxygen acceleration are similar to those of protons but appear in the larger oxygen diffusion region. In the exhaust, due to the geometry of the separatrices, the Hall electric field has $E_{x,GSM}$ and $E_{z,GSM}$ components. In their test particle simulation, these authors showed that an oxygen ion enters from the inflow region above the X-point and gains energy in the exhaust mainly from both the Hall electric field $E_{x,GSM}$ component and the reconnection electric field $E_{y,GSM}$. On the other hand, a proton mainly gains energy from the Hall electric field $E_{x,GSM}$ component. The Liu et al. [2015] study
gives an important picture of how the oxygen is accelerated and heated in the exhaust by the localized electric fields and further compares the results with spacecraft observations.

The oxygen ion sources involved in magnetotail reconnection include oxygen from the tail lobes which does not have a significant bulk velocity and pre-existing current sheet oxygen ions which have a drift velocity in the $Y_{GSM}$ direction caused by the meandering motion across the tail. The oxygen carries part of the cross tail current. The lobe oxygen may have a tailward flow but it is small compared with the drifting of the current sheet oxygen during reconnection within a thin current sheet. The pre-existing current sheet oxygen could result from earlier reconnection at other locations. As mentioned in Chapter 1, since the reconnection related electric fields are highly localized in association with reconnection-related structures, e.g., the exhaust, the separatrics, the Hall magnetic fields, the dipolarization fronts (DFs), and the outflow region downstream of the DFs, the oxygen acceleration is highly sensitive to its entry into these structures. Because the oxygen ions have much larger gyro-radii and a much larger diffusion region, the oxygen ion entry into the plasma sheet and possible non-adiabatic acceleration region should be much larger than that for protons. However, it is still unclear where the oxygen ions from the two sources enter the reconnection-related structures, what the oxygen entry-related acceleration and heating processes are, how sensitive these acceleration processes are to the entry into the plasma sheet and how the entry and acceleration processes differ between oxygen ions and protons.

In this chapter, by using 2.5D implicit particle-in-cell simulations (iPic3D), we investigate the behavior of oxygen ions entering the plasma sheet from the lobes and those pre-existing in the current sheet. We calculate the energy gain of these oxygen ions, and compare the results
with those for protons. These results will provide a comprehensive picture of oxygen ion acceleration and heating during magnetotail reconnection.

5.2. Simulation Set-up

In this chapter, we use the same simulation runs as in Chapter 4. Recall the simulations started with a Harris current sheet equilibrium in the X-Z plane without a guide field, i.e., the magnetic field is given by $B_x(z) = B_0 \tanh \left( \frac{z - z_0}{\Delta} \right)$ and the density by

$$n(z) = n_0 / \cosh \left( \frac{z - z_0}{\Delta} \right) + n_b,$$

where $B_0$ is the lobe magnetic field, $n_0$ is the plasma density at the center of the initial current sheet, $n_b$ is the background density (initially uniform in the simulation box), $\Delta$ is the half thickness of the initial current sheet, $z_0$ is the Z location of the center of the current sheet. Open boundary conditions are used. The simulation runs are the “O+ run”, the “equal-mass-density (EMD) run” and the “H+ run”. The “O+ run”, “EMD run” and “H+ run” included oxygen ion concentrations of 50%, 5% and 0%, respectively (the same concentration ratios are applied to both lobe and current sheet populations for each run). For convenience in this chapter, we refer to the three runs as “50% oxygen run”, “5% oxygen run” and “0% oxygen run”, which may correspond to the cases of the storm-time substorms, non-storm substorms with oxygen ions and non-storm substorms without oxygen ions, respectively. In Section 5.3-5.5, we only use the 50% oxygen run because it more easily illustrates the oxygen behavior for understanding the acceleration. The 5% oxygen run and 0% oxygen run will only be discussed with regard to the DF thickness in Section 5.6. If not specified, the results in this chapter are from the 50% oxygen run. The mass ratios of electrons, protons and oxygen with respect to the proton mass are $m_{e-} : m_{H^+} : m_{O^+} = 1/128 : 1 : 16$. The temperature ratios are $T_{e-} :$
In this chapter, in order to compare with spacecraft observations and other studies, the temperature for species ‘s’ is defined as \( T_s = \frac{1}{n_s} \text{trace}(\vec{P}_s) \), where \( \vec{P}_s \) is pressure tensor with

\[
\vec{P}_s = \int m_s (\vec{v} - \vec{V}_s)(\vec{v} - \vec{V}_s)f_s(\vec{r}, \vec{v}, t)d^3\vec{v} \quad \text{and} \quad \vec{V}_s = \int \vec{v}f_s(\vec{r}, \vec{v}, t)d^3\vec{v}
\]

and

\[
n_s = \int f_s(\vec{r}, \vec{v}, t)d^3\vec{v}
\]

The “heating” in this chapter refers to the increase in the temperature.

The normalization factors we used for the space and time scales are \( d_{H'} \) and \( \Omega_{H'} \), where

\[
d_{H'} = \left( \frac{n_0 e^2}{e_0 m_{H'}} \right)^{1/2} \quad \text{and} \quad \Omega_{H'} = \frac{eB_0}{m_{H'}},
\]

and \( c \) is the speed of light. The physical parameters used in this chapter are the pre-existing current sheet density \( n_0 = 0.2 \text{cm}^{-3} \), lobe magnetic field \( B_0 = 22nT \), and initial temperature \( T_{O+} = T_{H+} = 5T_e = 5keV \). Thus, the normalization factors are \( d_{H'} = 509.7 \text{km} \) and \( \Omega_{H'} = 2.1/s \), respectively. The simulation box has the dimensions \( L_x \times L_z = 300 \times 50 \times 50 \ d_{H'} \), which is equivalent to a region of about \( 23.5R_E \times 3.9R_E \) in the X-Z plane in the magnetotail. The X-point starts from a perturbation at the center of the box, i.e., at \( X_0 = 150 \) and \( Z_0 = 25 \). For all three species, there is a uniform background density \( n_{b,s} \) throughout the entire simulation box where \( n_{b,s} = 0.1n_{0,s} \) for each species, the subscript ‘s’ indicates different species, i.e., ‘e’, ‘H’ and ‘O’, \( n_0 = n_{0,e} = n_{0,H} + n_{0,O} \). The background populations mimic plasmas of lobe origin. More detailed set-up parameters for the simulation runs can be found in Chapter 4.

### 5.3. Ion Diffusion Region and Ion Energy Gain

In Sections 5.3 and 5.4, we will discuss the acceleration of oxygen of lobe origin. Generally, inside the ion diffusion regions, ions move non-adiabatically and are accelerated by electric fields. Outside the diffusion region, ions become frozen-in with the magnetic field and can experience adiabatic acceleration. Identifying the ion diffusion regions can help us analyze...
the ion acceleration processes. To identify the ion diffusion regions, we examine the perpendicular slippage, i.e., $|V_{\perp(H \text{ or } O)} - E \times B / B^2|$ [e.g., Goldman et al., 2015; Ashour-Abdalla et al., 2016]. Figures 5-1a and c shows the perpendicular slippage for the protons and the oxygen ions of lobe origin at time $\mathcal{O}t_H=61.24$, respectively. The slippage values are normalized to the local Alfvén speed for each ion species and are shown on a logarithm scale i.e., $\log_{10} (|V_{\perp(H \text{ or } O)} - E \times B / B^2| / V_{A(H \text{ or } O)})$. The reconnection occurs in the middle of the simulation box, at $(X_0,Z_0)=(150,25)$. Before we discuss the diffusion regions, we first define the regions we will focus on in this chapter, including the DFs, the exhaust, and the downstream outflow region. From examination of the $B_z$ component and the magnetic field lines in Figure 5-1, we define a DF as the leading edge of reconnected propagating flux at the center of the neutral sheet. The DFs are at around $|X-X_0|=24.6$ and $Z=25$ in Figure 5-1 (see also Figure 4-2 for a plot of the time evolution of the DFs). We define the exhaust as the outflow region with non-zero $B_z$, i.e., roughly $|X-X_0|<24.6$ in the outflow region in Figure 5-1. The DF is also the leading edge of the exhaust. The part of the outflow region downstream of the exhaust, i.e., roughly $|X-X_0|>24.6$, is called the downstream outflow region in this chapter. In the downstream outflow region, there is no reconnected flux $B_z$. There is current sheet plasma at $Z=25$ that is piled up by the advancing DFs. The locations of the DF, the exhaust and the downstream outflow region are shown in Figure 5-1d. The extent of the perpendicular slippage changes continuously in space as time advances for both of the protons and oxygen ions. To compare the spatial scales of the diffusion region between protons and oxygen, we plot thick black contours with the same value $|V_{\perp(H \text{ or } O)} - E \times B / B^2| / V_{A(H \text{ or } O)} = 0.3$ in both panels 1a and c. We did not choose $|V_{\perp(H \text{ or } O)} - E \times B / B^2| / V_{A(H \text{ or } O)} = 0$ as the identification criterion for the diffusion region because there always are some ions with extremely high energy in the velocity distribution tail. These ions
have larger gyro-radii than the spatial scale of the localized electric fields and thus lead to non-zero slippage. Selecting 0.2 or 0.4 or other values does not change the relative sizes between the proton and oxygen diffusion regions. By comparing the contours, we can see the oxygen diffusion region is much larger than the proton diffusion region. In panel a, the proton diffusion region roughly covers the region with $|X-X_0|<10$ and $|Z-Z_0|<3$ and the regions near the separatrices within about $|X-X_0|<14$. In Figure 5-1c the oxygen diffusion region covers the entire exhaust as well as part of the inflow region within $|Z-Z_0|<8$ and the regions near the separatrices extending far away from the X-point. Note that the perpendicular slippage near the separatrices could be a result of the diamagnetic effect. This cannot be interpreted as the diffusion of magnetic flux. As we will show in Section 5.4, near the separatrices, the oxygen ions are accelerated non-adiabatically by the Hall electric field and obtain more perpendicular kinetic energy. With larger gyro-radii, they start meandering in the downstream outflow region. The multi-scale ion diffusion regions in the presence of oxygen identified in Figure 5-1c are similar to those found by using the oxygen distribution function agyrotropy by Liu et al. [2015]. Note that in Figure 5-1, we only use ions of lobe origin to discuss the diffusion regions and energy gains. The current sheet species are mainly located in the region $24.5<Z<25.5$ and are pushed away from the X-point by the DFs. Thus, there are no current sheet components in the exhaust and outside of $24.5<Z<25.5$. This means that it is reasonable to discuss the diffusion regions and the energy gains around the X-point by using only the ion species of lobe origin.
Figure 5-1 (a) and (c): Perpendicular slippage for the protons and oxygen ions, respectively. They are normalized by the local Alfvén speed for each species and are shown on a logarithm scale. The thick black contours indicate a constant value $\log_{10}(0.3)$ for both ion species. (b) and (d): Energy gain $J^{(s)} \cdot E$ for the protons and oxygen ions, respectively, where ‘s’ indicates different ion species. The defined locations of the DF, the exhaust and the downstream outflow region are shown in Figure (d). The time of the snapshots is $\Omega_B t = 61.24$.

The ion bulk energy gain at different locations can be calculated from $J^{(s)} \cdot E$ in the rest frame, where current $J^{(s)}$ is calculated for ion species “s”. A positive $J^{(s)} \cdot E$ means particles are gaining energy from the electromagnetic fields and a negative $J^{(s)} \cdot E$ indicates ions are losing energy. Figures 5-1b and 5-1d show $J^{(s)} \cdot E$ for the protons and the oxygen of lobe origin, respectively. The protons gain energy in the exhaust, especially behind the DFs and near the separatrices within roughly $|X-X_0| < 26$. At the DFs, $J^{(s)} \cdot E$ for the protons shows a bipolar signature along X, which is mainly due to a bipolar electric field $E_x$ at the DF, which in turn is
caused by the divergence of the electron pressure tensor due to the pile-up of the electrons at
the DFs. The pile-up of electrons at the DF serves to neutralize the pile-up of the pre-existing
current sheet ions as discussed in Chapter 4. The protons neither gain nor lose energy in the
downstream outflow region. Compared with the protons, the oxygen ions gain energy in the
exhaust as well as in the separatrix region farther away from the X-point, i.e., the separatrices
in the downstream outflow region. The results of the ion energy gain show that near the
reconnection site, both protons and oxygen ions gain energy in their diffusion regions. This
indicates that the electric fields in these regions are important in accelerating the ions non-
adiabatically. These electric fields include the reconnection electric field $E_y$ near the neutral
sheet, $E_y$ due to the convection behind the DFs, the Hall electric field near the separatrices [e.g.,
Drake et al., 2008], and an ambipolar electric field $E_x$ near the neutral sheet due to the slower
motion along the outflow direction of oxygen ions compared to the electrons in the oxygen
diffusion region (see Chapter 4).
Figure 5-2 The (a) parallel (“||”) and (b) perpendicular (“⊥”) components (with respect to the magnetic field), and (c) X, (d) Y, (e) Z, components of the work \( J^{(H_+)} \cdot E \) for protons of lobe origin. The (f)-(j) panels show the same quantities and formats for the oxygen ions of lobe origin as in panels (a)-(c). The time of the snapshots is \( \Omega_{lt}=61.24 \).
In order to examine which electric field components accelerate the ions, Figure 5-2 shows the parallel (“||”) and perpendicular (“⊥”) components with respect to the local magnetic field, and X, Y, Z components of \( \mathbf{J}^{(H^+)} \cdot \mathbf{E} \) for the protons (Figures 5-2a to e) and oxygen ions (Figures 5-2f to j) of lobe origin. For the protons, the total energy gain (Figure 5-1b) is mainly from \( E_\perp \) (Figure 5-2b). Note that in Figure 5-2a, there are small regions at the DFs where protons gain energy from \( E_\parallel \). The \( E_\parallel \) at the DFs is generated by the parallel gradient of the electron pressure tensor, because there are peak values of \( P_{e,xx} \) and \( P_{e,zz} \) at the DFs due to the pile-up of electrons. Similarly, the perpendicular pressure gradient generates the bipolar electric field at the DFs and causes the bipolar \( \mathbf{J}_\perp^{(H^+)} \cdot \mathbf{E} \) in Figure 5-2b. As shown in Figure 5-1b, the proton energy gain at this time is mainly in the exhaust. From Figures 5-2c to e, we can see the total energy gain of protons (Figure 5-1b) is mainly from the \( E_x \) (Figure 5-2c), which is consistent with the results of Liu et al. [2015]. Near the separatrices, the protons gain energy due to the Hall electric field \( E_x \) and \( E_z \). Near the neutral sheet behind the DFs, they mainly gain energy due to the \( E_x \). It is worth noting that \( J_y^{(H^+)} \cdot E_y \) is negative in Figure 5-2d because the proton \( \mathbf{E} \times \mathbf{B} \) drift flow is opposite to the \( E_y \) outside the proton diffusion region. The negative \( J_y^{(H^+)} \cdot E_y \) does not mean protons are losing energy as it is canceled by the positive contributions from \( J_x^{(H^+)} \cdot E_x \) and \( J_z^{(H^+)} \cdot E_z \).

Similarly for the oxygen, the total energy gain (Figure 5-1d) is mainly from \( E_\perp \) (Figure 5-2g) as well. The oxygen ions do not show parallel acceleration at the DFs in Figure 5-2f, because unlike the protons, the oxygen outflow speed is small at the DFs due to their large inertia. In Figure 5-2g (also Figures 5-2h and j), near the separatrices, the oxygen ions gain energy due to the Hall electric field in both the exhaust and the downstream outflow region, which covers a much broader region than that of the protons. Near the neutral sheet in the exhaust (10<|X-
the oxygen ions gain energy mainly due to $E_y$ and $E_x$. In this region, $E_y$ is due to both the ion convection term and the Hall term, and $E_x$ is due to the ambipolar electric field in the oxygen diffusion region. *Liu et al.* [2015] have shown that protons and oxygen ions are accelerated by the Hall electric field and the reconnection electric field only in the exhaust. In the downstream outflow region, the Hall electric field, which is mainly along the $Z$ direction due to the geometry of the separatrices, significantly accelerates the oxygen ions toward the current sheet in front of the DFs. This has not been discussed before. Note that the oxygen diffusion region extends along the separatrices in the downstream outflow region as shown in Figure 5-1c.

### 5.4. Ion Entry and Acceleration

We focus on the acceleration of the lobe oxygen in this section and discuss the pre-existing current sheet oxygen in Section 5.5. In the previous section we showed (Figure 5-1b) that protons gain energy in the exhaust but do not gain or lose energy in the downstream outflow region. On the other hand, the oxygen ions (Figure 5-1d) gain energy in the exhaust, and also gain energy from the Hall electric field near the separatrices in the downstream outflow. Once the ions pass through the separatrices from the inflow region, they start to be influenced by the Hall electric field in the outflow region. In order to understand the ion acceleration, it is important to investigate the ion number flux entering the outflow region through different locations in the separatrix.

Figure 5-3a shows the $Z$ component of the proton number flux. Figure 5-3b shows the $X$ component (black) and the $Z$ component (green) of the proton number flux at $Z=30$ (horizontal magenta line in Figure 5-3a) as a function of $X$ location. The vertical dashed lines at $|X-X_0|=14$ show the locations of the separatrices. In this figure, the inflow region is $|X-X_0|<14$ and the
outflow region is $|X-X_0|>14$. At this $Z$ location, the proton flux in the inflow is mainly in the $Z$ direction and in the outflow is mainly in $X$ direction, especially in the downstream outflow region, e.g., $|X-X_0|>24.6$. 
Figure 5-3 The number flux of protons and oxygen ions of lobe origin is shown. (a) The color coded proton number flux in the Z direction. (b) The proton number flux in $X$ (black) and $Z$ (green) calculated at $Z=30$ along the horizontal magenta solid line in (a). (c) The $X$ (black), $Y$ (red), and $Z$ (green) components of the proton number flux as a function of $Z$ at $X=190$, along the vertical magenta dashed line in (a). The vertical dashed lines in both (b) and (c) indicate the locations of the separatrices. Panels (d)-(f) are plotted using the same formats as in panels (a)-(c), for the oxygen ions. The time of the snapshots is $\Omega_{Ht}=61.24$. 
We examined the proton number flux moving in the X (black), Y (red) and Z (green) directions as a function of Z location at X=190, as shown in Figure 5-3c. The proton flux was determined along the vertical magenta dashed line in Figure 5-3a. Note that this is in the downstream outflow region and at the moment when Figure 5-3c is taken, the propagating DF (at X=174.6) had not reached the location X=190. In Figure 5-3c, the two vertical dashed lines indicate the separatrices. The center of the current sheet is at Z₀=25. The proton flux in the X direction (black) shows a hump centered at Z=25 with a width comparable to the width of the outflow region at this location. The hump is formed by the outflowing protons which were previously accelerated in the exhaust and flow along the field to the downstream outflow region. This is similar to the results in the study by Eastwood et al. [2015], in which the authors demonstrate that the protons in the exhaust can overtake the propagating DF and populate the current sheet region downstream of the DF. The proton flux in Y shows two dips at Z=17 and Z=33 with widths of 2-3 d_H in Z. The locations and widths of the two dips are consistent with the locations and spatial scale of the Hall electric field E_z in this region (see discussion of Figure 5-4a). The two dips are proton E×B drifting populations due to E_z and B_x. The proton flux in the Z direction is almost zero compared to the other two components. Thus, the proton population in the downstream outflow region is mainly formed by the proton outflow in X direction previously accelerated in the exhaust. The protons entering from the separatrices in the downstream outflow region are limited to E×B drift in the region with the Hall electric field E_z.

Figures 5-3d to f are plotted for the oxygen ions in the same format as Figures 5-3a to c. Unlike the proton fluxes, the color-coded Z-direction flux in Figure 5-3d shows significant oxygen ion flux along Z from the separatrix to the current sheet in the downstream outflow.
region, i.e., \(|X-X_0|>24.6\). In Figure 5-3e, the oxygen ion flux at \(Z=30\) as a function of \(X\) location is different from that in Figure 5-3b for protons. In this figure, the oxygen number flux in the inflow region, i.e., \(|X-X_0|<14\) is in both \(X\) and \(Z\) direction. In the outflow region, oxygen number flux shows both the contribution of the \(X\) and \(Z\) components at \(14<|X-X_0|<24.6\). In the downstream outflow region, i.e., \(|X-X_0|>24.6\), the oxygen number flux is mainly in the \(Z\) direction, which is very different then for the protons in this region. We can see that in the downstream outflow region, the oxygen population comes from the \(Z\) direction, i.e., from the lobes, while the proton population comes from the \(X\) direction, i.e., from the exhaust. In Figure 5-3f the oxygen ion flux in the \(X\) direction is very small while the flux in the \(Z\) direction is significant. The oxygen flux in \(Z\) becomes large toward the current sheet due to the Hall electric field \(E_z\) after the oxygen ions pass through the separatrix. In Figure 5-4a, the Hall electric field \(E_z\) at \(X=190\) near the separatrix has a value of about \(20 \text{ mV/m}\) with a width of \(2-3 \text{ d}_H\) in \(Z\). For the initial \(T_{O+}=5\text{keV}\) oxygen in the \(22nT\) lobe \(B\) field, the gyro-radius is about \(1313 \text{ km}\), i.e., about \(2.6 \text{ d}_H\). The gyro-radius is comparable to the spatial scale of \(E_z\) in the \(Z\) direction, which means that the oxygen ions in that area experience non-adiabatic acceleration along the \(Z\) direction. This is very different behavior than that seen for protons, which experience \(E \times B\) drift in the \(Y\) direction (shown the proton number flux in \(Y\) direction (red) in Figure 5-3c) when passing through the separatrices.

Now, let us estimate the oxygen energy gain due to the \(E_z\) component. The work done by a \(20\text{ mV/m}\) \(E_z\) over a distance of \(2.5\text{d}_H\) in \(Z\) is about \(4.0 \times 10^{-15} \text{ J}\). This work is equivalent to the kinetic energy of an oxygen ion with a speed of \(547 \text{ km/s}\), and the gyro-radius of this oxygen ion in the \(22nT\) \(B\) field is about \(4.15\times10^3 \text{ km}\), i.e., about \(8.3 \text{ d}_H\). This means that after being accelerated by \(E_z\), the oxygen gyro-radii become much larger and the oxygen ions are able to
reach and even pass the center of the current sheet. Note that in the downstream outflow region, the magnetic field is dominated by $B_x$ and its strength does not change by much during the evolving simulation. The above calculation of the gyro-radius in a constant B field is still a good estimate. Also note that the time scale for the simulation is short compared to that for oxygen motion, as it only allows the oxygen ions to finish about half a gyro-orbit after entering the outflow region. The significant negative flux in $Y$ (red) in Figure 5-3f indicates an increased oxygen gyro-velocity in the $Y$ direction. This will be confirmed by tracking single particle motion in Section 5.6. In general, these accelerated oxygen ions have higher temperatures than their initial state and populate the downstream outflow region even before the DF reaches them.

Figure 5-4b shows the oxygen phase space density as a function of $V_z$ and $Z$ at $X=190$ which is highlighted along the green dashed line in Figure 5-4a. Figure 5-4c shows phase space density for protons using the same format as Figure 5-4b. Figure 5-4d shows the electric field components as a function of $Z$ at $X=190$. The vertical dashed lines show the location of the separatrices. From Figures 5-4b and 5-4d, we see that after passing through the separatrices, the oxygen ions are accelerated toward the current sheet by $E_z$. The average speed after being accelerated is about 500 km/s, which is close to our estimate of 547 km/s determined from the work done by $E_z$. After acceleration, the oxygen ions start to gyrate with much larger gyro-radii and form a population hotter than the initial one in the downstream outflow region. Such a non-adiabatic acceleration process is not evident in Figure 5-4c for the protons, which experience adiabatic motion at the same locations due to their much smaller gyro-radii.
Figure 5-4 (a) The $E_z$ component as a function of $X$ at $\Omega_{Ht}=61.24$. (b) The oxygen phase space density as a function of $V_z$ and $Z$ at $X=190$, i.e., along the green dashed line in (a). (c) Proton phase space density in the same format as (b). (d) Electric field components along the green dashed line in (a). The black vertical dashed lines highlight the locations of the separatrices. (e) The evolution of $E_z$ at $X=168$, i.e., along the magenta dashed line in (a), as a function of $Z$. The time is in units of the inverse proton gyro-frequency.

The $E_z$ component of the Hall electric field works as a potential well in the $Z$ direction. It is good to compare the oxygen non-adiabatic acceleration by $E_z$ in the downstream outflow region with that in the exhaust. During the time-evolution of the reconnection, this $E_z$ forms a time-dependent potential well. By comparing the time scale of the oxygen motion and the changing rate of the magnitude and the width of the potential well, we can identify different energization effects. Figure 5-4e shows the $E_z$ at $X=168$ (indicated by vertical magenta dashed line in Figure 5-4a) as a function of $Z$ and time (normalized to the inversed proton gyro frequency $\Omega_{H^{-1}}$). At time $\Omega_{Ht}=50$, the DF passes this $X$ location (i.e., $X=168$). At times $20<\Omega_{Ht}<50$, before the DF passes $X=168$, Figure 5-4e shows $E_z$ in the downstream outflow region. In this interval, both the magnitude of $E_z$ and the width of the potential well change on the time scale of oxygen gyro-motion (note that the oxygen gyro frequency $\Omega_O=16\Omega_H$). From $\Omega_{Ht}=20$ to $\Omega_{Ht}=50$, it is about $2\Omega_O^{-1}$ and the width of the potential well changes about $12d_H$. The oxygen ions accelerated on one side of the peak of the $E_z$ do not necessarily reach the other side of the peak of the $E_z$ at a later time. They tend to stay around the neutral sheet with meandering motion while the downstream outflow region expands in the $Z$ direction. At times $\Omega_{Ht}>50$, Figure 5-4e shows $E_z$ in the exhaust. From $\Omega_{Ht}=50$ to $\Omega_{Ht}=95$, it is about $3\Omega_O^{-1}$ and the width of the potential well changes about $4d_H$. Therefore, neither the magnitude of $E_z$ nor
the width of the potential well has significantly changed compared to the downstream outflow region. This is what you would expect for a nearly static potential well. Therefore, for the oxygen acceleration in the exhaust, the effect of the $E_z$ contributes to bounce motion rather than energization. Note that we have compared the Z-T plots of $E_z$ at different X locations (not shown) and found that (1) the evolution is related to the outflow flux transport instead of the temporal change of the reconnection rate; (2) the expansion rate of the potential well in the downstream outflow region is identical at different locations, which means that the fast expansion is an unique property of $E_z$ potential well in the downstream outflow region.

In summary, the acceleration processes for the lobe protons and oxygen ions are very different during magnetotail reconnection. The oxygen ions entering both the exhaust and the downstream outflow region gain energy, while the protons can experience significant acceleration only when they enter the exhaust. This is consistent with the energy gain results in Figure 5-1b and d. The protons can be accelerated in the exhaust along the outflow direction. The accelerated protons flow along the field lines and further populate the downstream outflow region. A portion of the oxygen ions enters the exhaust and gains energy from the Hall electric field, reconnection electric field and the ambipolar electric field. Another portion enters the downstream outflow region directly from the lobe and is accelerated by the Hall electric field $E_z$.

5.5. Oxygen Phase Space Density in front of the DF

In Sections 5.3 and 5.4, we discussed the oxygen from the lobe source that enters the separatrices, the exhaust, and the downstream outflow region. In this section, we discuss the oxygen ions pre-existing in the current sheet along with the lobe source and their relationship to DF regions.
Figure 5-5a shows the magnetic field components at the center of the current sheet ($Z_0=25$) as a function of $X$ location. The peak in $B_z$ indicates a DF and is at $X=174.6$. Note that the X-point is at $X_0=150$ and the DF is propagating in the positive $X$ direction. The region at $X<174.6$ is in the exhaust. As described before, the DF is the leading edge of the reconnected flux pile-up region. If we look from the outflow region toward the X-point, we can see that $B_z$ starts to increase at $X=177.2$ and reaches a peak at $X=174.6$. We define the region $174.6<X<177.2$ as a DF layer. The pre-existing current sheet population exists to the right of the DF, i.e., $X>174.6$. Figure 5-5b shows the oxygen phase space density as a function of $V_x$ and $X$ at the center of the current sheet ($Z_0=25$). In Figure 5-5b, the region downstream of the DF for $174.6<X<196$ contains both reflected oxygen ions and current sheet oxygen ions. For $X>196$, there are only pre-existing current sheet oxygen ions. Positive $V_x$ is away from the X-point. The black dashed line indicates the location of the DF. In Figure 5-5b, the pre-existing current sheet oxygen distribution is shown as a beam centered at $V_x=0$ in the region with $X>177$. Initially, the oxygen ions in the pre-existing current sheet have a bulk velocity $V_y=1787$ km/s (shown in Figure 5-5d), which essentially forms one ion current component of the Harris current sheet [e.g., Galeev, 1982]. In order to facilitate reconnection in the simulation, we used a high value for the cross-tail bulk velocity even though it is higher than observed [e.g., Kistler et al., 2005]. But this high $V_y$ value does not change our conclusions. The average initial $V_x$ is zero. Another beam centered at about $V_x=1.4\times10^3$ km/s for $177<X<195$ consists of the reflected current sheet oxygen ions. Note that the center $V_x$ of this beam decreases slightly with increasing X locations. At this moment, the speed of the DF is about $V_{DF}=735.4$ km/s. In the frame of the moving DF, the pre-existing current sheet oxygen ions have an initial $V_y$ and a $V_x=-V_{DF}$ toward the DF. They deviate and are reflected back with $V_x=+V_{DF}$ due to the $B_z$ of the DF. In the rest frame,
the $X$ component of the bulk velocity of these oxygen ions increases from zero to $2V_{DF}$. They gain energy from the induced electric field of the changing $B_z$ in the rest frame. This mechanism was discussed previously by Zhou et al. [2010, 2012, 2015] based on observed DF events and test particle modeling.
Figure 5-5 (a) The $B$ field at the center of the current sheet ($Z=25$). The black vertical dashed line indicates the location of the DF. (b) Oxygen phase space density as a function of $x$ and $V_x$, i.e., $f(x, V_x)$, at $Z=25$. The other velocity components are integrated. (c) – (d): The oxygen phase space density in the $Vx-Vz$ and $Vy-Vz$ planes, i.e., $f(V_x, V_z)$ and $f(V_x, V_y)$, at $X=178.5$ and $Z=25$, which are highlighted by a green dashed line in (b). In (d), six peaks in the distribution function are numbered. Peaks 1, 2, 3 consist of pre-existing current sheet O$^+$ ions. Peaks 4, 5, 6 are formed by background (lobe) O$^+$ ions. (e) The proton phase space density in the same format as (b). The black vertical dashed line indicates the location of the DF. Note that the Y range of (e) is different than (b).

In this study, we found a new feature in the phase space density of the reflected current sheet oxygen. In Figure 5-5b, at around $X=177.5$, there is a third population in addition to the current sheet and reflected oxygen ions centered at about $V_x = 750$ km/s between the initial current sheet population (centered at $V_x=0$ km/s) and the reflected beam (centered at about $V_x=1.4\times10^3$ km/s). The center $V_x$ of this third beam increases with $X$ and connects to the reflected beam at about $X=192$. The third beam is formed by oxygen reflected at earlier times by the propagating DF. In the simulation, DFs form at the beginning of the fast reconnection phase, and their speed increases as they propagate away from the X-point (see Figure 4-3). As they propagate, they continuously reflect the pre-existing current sheet oxygen ions. These reflected oxygen ions gain bulk velocity $2V_{DF}$ in $X$, and they actually record the instantaneous speed of the DF during early times. Because there is no magnetic field in the plane perpendicular to $X$ in front of the DFs, the reflected oxygen ions will retain their speed in $X$. As the speed of the DF increases with time, it will be able to catch up to the oxygen ions that were reflected earlier. The third beam at $X=178.5$ is the reflected oxygen that is about to be overtaken by the
DF. In other words, the third beam is actually a part of a reflected beam with a hook shape, which records the speed history of the DF. Note that the hook shape in the reflected beam can only appear clearly when the width of the beams is small so that they do not overlap each other.

For a typical ion temperature of 5 keV in the magnetotail plasma sheet, the 5-keV oxygen ions have a thermal velocity $V_{th,O}=173 \text{ km/s}$, while the 5-keV protons have thermal velocity $V_{th,H}=692 \text{ km/s}$, which is comparable to the speed of some DFs. Figure 5-5e shows the proton phase space density with the same format as Figure 5-5b. From Figure 5-5e, we can see that even though there are proton reflection beams, the incident and reflected beams overlap each other because their typically thermal velocities are large when compared to their bulk velocities. Therefore, we expect that it would be more difficult for spacecraft to observe a hook shape in the reflected protons than it is in the reflected oxygen. This feature of oxygen phase space density can be applied to determine the speed history of DFs observed by spacecraft.

Figures 5-5c and 5-5d show the oxygen phase space density at the center of the current sheet at $X=178.5$ projected in the $V_x - V_z$ and $V_x - V_y$ planes, respectively. The locations at which the distribution functions were calculated are highlighted by a green dashed line in Figure 5-5b. The current sheet and reflected oxygen beams in Figure 5-5b are also found as peaks centered at different $V_x$ in Figures 5-5c and 5-5d. In Figure 5-5c, the distribution functions are almost symmetric in the Z direction and the reflection beams are shown in the positive X direction only. In Figure 5-5d, there are six features, which are marked by the numbers 1-6. Numbers 1-3 are formed by the pre-existing current sheet oxygen ions, while the numbers 4-6 are formed from background (lobe) oxygen ions. Features 1 and 5 correspond to the initial current sheet and background oxygen ions, respectively. Features 2 and 3 are reflected current sheet oxygen ions as discussed above. Feature 4 involves some lobe oxygen ions that entered
the exhaust and were pushed by the moving DF and/or the Hall electric field $E_x$ component. Finally, the peak at 6 is formed by the lobe oxygen ions that entered the downstream outflow region and were accelerated by the Hall electric field $E_z$.

Three examples of oxygen ion trajectories taken from the simulation output that gave peaks 2-4 are shown in Figures 5-6a to c, respectively. The trajectories are shown by gray curves. They are plotted on top of the $B_z$ (Figures 5-6a and 5-6b) and $E_x$ (Figure 5-6c) values. The trajectories show the positions of these oxygen ions starting from $\Omega_{Ht}=21.05$ and continuing until they reach $X=178.5$ and $Z=25$ at $\Omega_{Ht}=61.24$, the moment at which the phase space density results are captured in Figure 5-5. The times for each snapshot are the times when these oxygen ions encountered the DF. Each green filled circle represents the oxygen ion location at the time of its DF encounter. The field line passing through the location of the particle is highlighted in green.

The particle in Figure 5-6a, is a pre-existing current sheet oxygen ion that was in the distribution and labeled 2 in Figure 5-5d. The $qV_O \times B$ and $qE$ forces that occur when the ion interacts with the DF are indicated by blue and red lines, respectively, starting from the center of the ion and pointing in the direction of the forces. The lengths of the lines are proportional to the magnitudes of the forces. This oxygen ion is reflected ahead of the DF at $\Omega_{Ht}=32.5$. The $qV_O \times B$ term in the Lorentz equation is dominant along $X$. This results from the initial $V_y$ of a current sheet oxygen ion combined with the strong $B_z$ of the DF.
Figure 5-6 (a)–(c): The Peak 2 particle trajectory (gray) overlaying $B_z$ at $\Omega_{Bt}=32.5$, Peak 3 particle trajectory (gray) overlaying $B_z$ at $\Omega_{Bt}=47.8$, and Peak 4 particle trajectory (gray) overlaying $E_x$ at $\Omega_{Bt}=30.6$. The corresponding particles at the times specified are plotted as filled circles (green) near (a) X=160 (b) X=167 and (c) X=158. The field lines on which each particle is located are indicated by the green curves. The $qV \times B$ and $qE$ forces on the particles are indicated as blue and red lines, respectively. The directions of the forces are indicated by the blue and red lines pointing in the direction away from the center of the particles. The lengths of the blue and red lines are proportional to the magnitudes of the forces. In (a) and (b), the Peak 2 and Peak 3 particles encounter the DF in succession due to the $qV \times B$ force. In (c), the Peak 4 particles encounter the DF due to the electric field $E_x$ at the DF.

Figure 5-6b shows an oxygen ion from the structure marked 3. The format is the same as in Figure 5-6a. The DF reaches this ion at $\Omega_{Bt}=47.8$ and the $qV_O \times B$ force is the dominant term in the Lorentz equation. Ions in peak 3 were reflected by the DF later than those in peak 2. The trajectory of this ion also shows meandering motion and trapping around the current sheet, which is projected onto the X-Z plane.

Figure 5-6c shows a lobe oxygen ion contributing to peak 4. Here the background is the electric field $E_x$ component. This oxygen ion entered the exhaust from the northern lobe. It interacted with the DF at about $\Omega_{Bt}=30.62$. The $qE$ term is dominant in this case because the lobe oxygen ion does not have a significant initial $V_y$, unlike current sheet ions. Note that although the trajectory leads to the other side of the inflow region in this snapshot, the outflow region will expand in $Z$ later in time and the oxygen ion will bounce in the outflow region due to the Hall electric field $E_z$ without crossing the separatrix on the other side. This will be discussed further in Section 5.6. It is worth noting that this oxygen ion also obtains $V_x$ due to a
push by the Hall electric field $E_x$ near the separatrix when it bounces in the exhaust. The $V_x$ gain due to such a push is strongly dependent on its location near the separatrix. In addition, this lobe oxygen ion gains a $Y$ component ($V_y$) due to the effect of the $E_y$ in the exhaust. The gain in $V_y$ also is strongly dependent on its location and the period of time it stays in the exhaust. For lobe oxygen ions with different initial states, i.e., different locations and velocity when they enter the exhaust, the gain in both $V_x$ and $V_y$ may be very different due to the different possible encounters with the electric fields in the exhaust. This explains why peak 4 is much more diffusive in the $V_x - V_y$ plane than the other features.

In general, the pre-existing current sheet oxygen is accelerated by reflection off of a DF. These oxygen ions form a hook-shaped reflected beam in the phase space ($X-V_x$) due to their continuous reflection as the DF increases speed. This hook-shaped reflected beam records the speed history of the DF. However, this feature is hard to observe for protons, because the typical thermal velocity of protons in the plasma sheet is comparable to the speed of the DF, so the proton reflection beams are too wide and overlap each other. In addition there are multiple peaks in $V_x - V_y$ phase space in front of the DF due to the different sources of oxygen ions.

5.6. Discussion

As shown in Section 5.3 and similar to Liu et al. [2015], we found that in the exhaust, the $E_x$ and the $E_y$ produce the main oxygen ion energy gain. In addition, we found that in the downstream outflow region, the Hall electric field $E_z$ near the separatrix produces the main energy gain. This indicates that the entry locations on the separatrix of the lobe oxygen ions actually serve to determine the acceleration processes affecting them. In Section 5.4, we showed that besides entering the exhaust near the X-point, there are significant numbers of lobe oxygen ions that pass through the separatrix in the downstream outflow region and are accelerated by
the Hall electric field $E_z$. Because Hall fields propagate much faster than DFs [Shay et al., 2011, Lapenta et al., 2013], they can range much further downstream in the outflow region than DFs and therefore accelerate a significant number of oxygen ions before the DFs reach those same regions. To further discuss the oxygen ion energy gain due to different entry locations, we look at two example oxygen ions and determine their energy gains as a function of time.

In Figure 5-7, an oxygen ion is tracked during the PIC simulation. This is the same ion represented in Figure 5-6c. Figure 5-7c shows several physical parameters along its trajectory as a function of time. Panels 7a and b show the trajectory superimposed on field lines and $E_x$ during two times. This oxygen ion reached the outflow region by entering the exhaust from a location close to the X-point. Panels 7a and 7b show its positions at the times when it reached the northern and southern separatrices, respectively. These two times are indicated by two vertical dashed lines in Figure 5-7c. After entering the exhaust, this ion was influenced by the electric fields in the exhaust. After passing through the northern separatrix, it encountered a DF before crossing the current sheet. In panel 7b it was deflected back into the downstream current sheet by the Hall electric field. Here, the electric field force term in $Z$ is larger than the Lorentz force at the time it was deflected. This indicates the electric potential well had a large contribution, trapping this oxygen ion in the outflow region in addition to its meandering motion. Such an effect of the electric potential well was discussed in previous studies [Wygant et al., 2005; Aunai et al., 2011]. Finally, this oxygen ion reached a location in front of the DF and contributed to structure 4 in the phase space density in Figure 5-5d. It gained about $45keV$ energy through this process. As was found by Liu et al. [2015], $E_x$ and $E_y$ in the exhaust equally provided the main work done on this ion, while $E_z$ mainly helped it bounce without providing a significant energy gain.
Figure 5-7 Tracking one oxygen trajectory (gray) overlaying $E_x$ at (a) $\Omega_{Ht}=26.8$ and at (b) $\Omega_{Ht}=42.1$. The field lines on which this oxygen ion is located are indicated by the green curves. (c) The time evolution of X, Z, velocity, B field, E field, kinetic energy and energy gain of this oxygen ion. The two vertical lines indicate the times during which (a) and (b) are plotted.

Figure 5-8 shows another oxygen ion which entered the downstream outflow region without being in the reconnection exhaust. This one is from feature 6 in Figure 5-5d. This ion initially was in the northern lobe (Figure 5-8a). It crossed the expanding separatrix at $\Omega_{Ht}=50$. As it crossed the separatrix it was energized. At that moment, this oxygen ion is at about $X=180$, while the DF is at about $X=169$. The oxygen was far away from the exhaust and then further accelerated to the center of the current sheet. Note that it had completed only about half of a gyro-orbit after entering the separatrix. Even though this oxygen ion did not enter the exhaust, it gained about $30 \text{ keV}$ energy from this process - due to the Hall electric field.
Figure 5-8 (a) Tracking one oxygen trajectory (gray) overlaying $E_z$ at $\Omega_H t=49.8$. The field lines on which this oxygen ion is located are indicated by the green curves. (b) The time evolution of X, Z, velocity, B field, E field, kinetic energy and energy gain for this oxygen ion. The vertical line indicates the times at which (a) is plotted.

In Figures 5-9a and 5-9b we have plotted the proton temperature for the lobe component and the pre-existing current sheet component, respectively. The demonstration of temperature provides a reference for comparison with spacecraft observations. Recall that the DFs are at $|X-X_0|=24.6$. We call the region $|X-X_0|<24.6$ the exhaust and the region $|X-X_0|>24.6$ the downstream outflow region. The initial temperature was 5 keV. We found that almost all the heated lobe protons populate the outflow region or the region near the X-point. The protons are strongly heated near the X-point due to the reconnection electric field $E_y$ during their meandering motion [e.g., Moses et al., 1993]. In the exhaust, near the X-point, the high temperature results from the counter-streaming distributions of the protons due to the meandering motion and the Hall electric fields between separatrices [e.g., Hoshino et al., 1998; Wygant et al., 2005; Drake et al., 2009]. In the region of the exhaust further away from the X-point, there is also acceleration that occurred when the ions were picked up by the convecting magnetic field [Drake et al., 2009]. In the downstream outflow region, the high temperature is mainly in the parallel direction with respect to the magnetic field (not shown) and it is formed by the heated counter-streaming protons flowing along the magnetic field lines from the exhaust, which is consistent with the dominant outflow flux $n^{(H^+)}V^{(H^+)}_z$ shown in Figure 5-3c. Note that in the downstream outflow region since there are few inflow protons and the outflow convecting $B_z$ has not reached this region there is little acceleration in the pick-up process. Therefore, in the downstream outflow region, there is no locally-generated counter-streaming inflow or pick-
up ion signature like we found in the exhaust. In Figure 5-9b, the high temperature of the current sheet protons is due to the combined effects of the reflected beam and the pre-existing current sheet population [e.g., Zhou et al., 2010]. Note that in Figures 5-9b and 5-9d, to minimize noise, the grids with the pre-existing current sheet ion density lower than 0.002 cm$^{-3}$ were ignored. Proton heating in the exhaust in the presence of oxygen ions is weaker than that for the case without oxygen, because the oxygen reduces the reconnection rate and thereby the reconnection electric field. The peak energy in Figure 5-9b is about 13 keV while in the simulation with 0% oxygen it is about 23 keV.
Figure 5-9 Proton temperature for the (a) lobe component and (b) pre-existing current sheet component. Oxygen temperature for the (c) lobe component and (d) pre-existing current sheet component. Note that in order to avoid the noise due to insufficient particles in a cell, the grids with a density lower than 0.002 cm$^{-3}$ were ignored in (b) and (d). The time of the snapshots is $\Omega_{H,t}=61.24$.

Figures 5-9c and 5-9d show the oxygen temperature for the lobe component and the pre-existing current sheet component, respectively. The oxygen ions also are counter-streaming in the exhaust, which is consistent with oxygen ion observations by Wygant et al. [2005]. Because the oxygen diffusion region covers the entire exhaust and broader regions near the separatrices, the oxygen ions are accelerated by the electric fields not only near the X-point but also in a much broader region farther away from the X-point. In the downstream outflow region, the oxygen ions are accelerated by the Hall electric field $E_z$ near the separatrices and then meander around the neutral sheet during the extension of the downstream outflow region. They contribute to the high temperature in this region. It is worth noting that unlike the protons, as shown in Figure 5-3f, the oxygen ion outflow flux $n^{(O^+)}V^{(O^+)}$ is very small, which indicates that the heated oxygen in the exhaust has not yet propagated to the downstream outflow region due to the large inertia of the oxygen ions. The only source for the hot population in the downstream outflow is the oxygen directly entering this region through the separatrices. Similar to the protons, the reflected oxygen beam also contributes to the high temperature of the pre-existing current sheet component in front of DFs.

The DF is believed to be an important magnetic structure frequently related to an ion dissipation region [e.g., Zhou et al., 2010; Drake et al., 2015]. There are also Hall currents [e.g., Fu et al., 2012] and plasma waves associated with DFs [e.g., Sergeev et al., 2009; Zhou et al.,
2009, 2014; Deng et al., 2010]. The spatial scales of DFs determine whether the electrons and ions are frozen-in within this layer. In spacecraft observations, the thicknesses of DFs in the near-Earth magnetotail are on the scale of the ion inertial length or ion gyro-radius [e.g., Runov et al., 2009], so a DF layer is believed to be a region where ions are not frozen-in. To study how the scale of this ion diffusion region, i.e., the thickness of the DF layer, is influenced by the oxygen ion content during different magnetospheric activity levels, we compare the DFs in the 50% oxygen run, the 5% oxygen run, and the proton (0% oxygen) run. Except for the concentration of oxygen ions, the three runs use the same parameters. Note that the current sheet thicknesses are the same for the three runs even though the oxygen concentration is different. The current sheet thickness is independent of the oxygen content as suggested in both of the theory of Harris current sheets (i.e., the current sheet thickness is determined by the ambient magnetic field, the plasma temperature and the cross-tail speed) [e.g., Galeev, 1982] and observations [e.g., Liu et al., 2014]. Figure 5-10 shows the DF profiles at the center of the current sheet. For the 50% oxygen run (Figure 5-10a) the DF layer, as indicated by the vertical dashed lines, has a thickness \(2.6 \, d_H\). In the DF layer, as shown in Figure 5-10b, the pre-existing current sheet ions (blue and red lines) pile up and have a much greater number density than the background (dashed lines) ions. Figures 5-10c and 5-10d, for the 5% oxygen run, and Figures 5-10e to f, for the proton run, show that the DF thickness is \(1.4 \, d_H\) and \(1.2 \, d_H\), respectively. The pre-existing current sheet ions also pile up at the DF layer during these two runs. For all the runs, the DFs separate the low density plasma and the high density plasma in the neutral sheet, as previously observed [e.g., Runov et al., 2009].
50% oxygen run

$B_z$ [nT]

$N_e$ (cs)
$N_i$ (cs)
$N_e$ (b)
$N_i$ (b)

5% oxygen run

$B_z$ [nT]

$N_e$ (cs)
$N_i$ (cs)
$N_e$ (b)
$N_i$ (b)

0% oxygen run

$B_z$ [nT]

$N_e$ (cs)
$N_i$ (cs)
$N_e$ (b)
$N_i$ (b)
Figure 5-10 (a) The $B_z$ field along X. (b) The current sheet density (‘cs’, indicated as solid lines) and background density (‘b’, indicated as dashed lines) for electrons (‘e’) and protons (‘H’) and oxygen (‘O’), for the 50% oxygen run. (c)-(d) and (e)-(f) are for the 5% oxygen run and the 0% oxygen run, respectively, using the same formats as (a)-(b). All of these quantities are measured at the center of the current sheet along the X-axis. The DF layer is highlighted by the double vertical dotted lines for all the runs. The thicknesses of the DF's are shown between the dotted lines.

Because of the pressure balance across a DF, calculating its thickness is equivalent to calculating the penetration depth of the pre-existing current sheet ions. Thus the dynamics of the pre-existing current sheet ions in a DF determine its thickness. As we showed in Section 5.5, when a DF hits the pre-existing current sheet ions, the major force exerted on them is the Lorentz force which is due to their significant bulk velocity in the Y direction. For each of the pre-existing current sheet ions penetrating into a DF with mass $m$, the Lorentz force $f$ in the X direction changes the momentum in a time interval $\Delta t$, i.e., $f_x \Delta t = m \Delta v_x$, where $\Delta v_x$ is the change of its speed in the X direction. When considering all the penetrating ions, we use an averaged form $<f_x/m> \Delta t^* = \Delta v_x^*$, where $<f_x/m>$, $\Delta t^*$ and $\Delta v_x^*$ are averaged over all the protons and oxygen ions. The $<f_x/m>$ term can be obtained from the simulation data. Assuming that $\Delta t^*$ is the entire interval needed to finish the reflection, we can then estimate $\Delta v_x^* \approx 2V_{DF}$, where $V_{DF}$ is the speed of the DF. We assume $\Delta t^*$ to be $c_0D/V_{DF}$, where $c_0$ is a constant and $D$ is the penetration depth. If $c_0=1$ when the DF propagates a distance $D$, then all the ions within a distance $D$ are reflected away from the DF. In this case, there will not be any pile-up of ions in the DF, because the number of reflected ions is the same as the number of the ions flowing in the DF frame. In Figures 5-10b, d and f, we can see pile-up ions in the DF for all three runs. Note that the initial
total number density of the pre-existing current sheet ions is $0.2 \text{ cm}^{-3}$ in the 50% oxygen ion run, $0.1 \text{ cm}^{-3}$ for protons and $0.1 \text{ cm}^{-3}$ for oxygen ions. In the 5% oxygen run, the density is $0.19 \text{ cm}^{-3}$ for protons and $0.01 \text{ cm}^{-3}$ for oxygen. For all three runs, the pile-up ion densities are about twice as large as the initial values. Thus we can consider $c_0=2$ to be a good approximation, i.e., $\Delta t^*=2D/V_{DF}$. Note that for the 50% oxygen run, the pile-up oxygen density is about three times as large as its initial value, while for the protons it is about 1-2 times, so using “twice” the initial value is an acceptable estimate. Finally, we can estimate the penetration depth, i.e., $D \approx V_{DF}^2/(\langle f/M \rangle)$. The results for the three runs are shown in Table 5-1. They are very close to the DF thickness we measured from the $B_z$ profile. Note that as shown in Chapter 4, the speed of the DF is strongly influenced by the inertia of the pre-existing current sheet, i.e., the concentration of oxygen ions in the current sheet.

Table 5-1: Calculation of the penetration depths of the pre-existing current sheet ions for the three runs and comparison with the respective dipolarization front thicknesses.

<table>
<thead>
<tr>
<th>O+ concentration</th>
<th>$V_{DF}$ (km/s)</th>
<th>$\langle f/M \rangle$ (km/s²)</th>
<th>$D=V_{DF}^2/(\langle f/M \rangle)$ (Note: $1,d_H \approx 509.7, \text{km}$)</th>
<th>DF thickness ($d_H$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>50%</td>
<td>735.4</td>
<td>433.1</td>
<td>$1248.7,\text{km} \approx 2.45,d_H$</td>
<td>2.6</td>
</tr>
<tr>
<td>5%</td>
<td>1435.8</td>
<td>2643.5</td>
<td>$779.8,\text{km} \approx 1.53,d_H$</td>
<td>1.4</td>
</tr>
<tr>
<td>0%</td>
<td>1470.8</td>
<td>3143.8</td>
<td>$688.1,\text{km} \approx 1.35,d_H$</td>
<td>1.2</td>
</tr>
</tbody>
</table>
5.7. Conclusion

In this chapter, we studied the oxygen heating process during magnetotail reconnection. We used a 2.5D implicit simulation code, iPc3D, to investigate magnetotail reconnection in the presence of an oxygen ion concentration of 50%, similar to what might be found in a storm-time substorm. The simulation starts with a Harris current sheet without a guide field. In order to understand the entire picture of the oxygen heating in magnetotail reconnection, we discuss the lobe oxygen ions and the pre-existing current sheet oxygen ions individually by considering the sources: the lobe oxygen ions and the pre-existing current sheet oxygen ions.

For the lobe ions, the acceleration due to the electric field around the reconnection regions is closely related to the locations where they enter the outflow region through the separatrices. The electric fields in the outflow region mainly contribute to the ion acceleration. In the exhaust, \textit{i.e.}, the outflow region between the X-point and the DFs, the electric fields include the Hall electric field near the separatrix, the reconnection electric field, and an ambipolar electric field due to the differential speeds along the outflow direction between the oxygen ions and the electrons (see Chapter 4). In the downstream outflow region, \textit{i.e.}, the outflow region downstream of a propagating DF, the electric field is mainly the Hall electric field near the separatrices. For the lobe protons, a fraction of them enter the exhaust and gain energy mainly from $E_x$, while another fraction enters the downstream outflow region and begins $E \times B$ drift along the Y direction near the separatrices without gaining much energy. The heated proton population in the downstream outflow region consists of the heated protons in the exhaust which flow downstream along the field lines. For the lobe oxygen ions, a fraction of them entering the exhaust gain energy from the electric fields mentioned above. Unlike the protons, the rest of the lobe oxygen ions directly enter downstream of the outflow region and are accelerated by
the Hall electric field. With this acceleration, they are able to reach and meander around the center of the current sheet and become a heated oxygen component populating the downstream outflow region in front of the DFs without entering the exhaust. Because of their large inertia few heated oxygen ions are found in the exhaust flowing downstream in the outflow region. This is in sharp contrast to the protons. These results were confirmed by either tracing oxygen motion in the PIC simulation or calculating single oxygen ion trajectories in the electric and magnetic fields from the PIC simulation (not shown).

For the pre-existing current sheet oxygen ions, which form the initial current sheet and have a significant bulk velocity in the cross-tail direction ($Y_{GSM}$ direction), we mainly focused on the ion heating related to DFs. By examining the oxygen phase space density as a function of $V_x$ and $X$, we demonstrated that there is an oxygen reflected beam with a hook shape in front of the DF. This hook-shaped reflected beam results from the increasing speed of the DF and can be used to track the DF’s speed history during the simulation. This hook feature is difficult to identify for protons, because the typical proton thermal velocity in the tail plasma sheet is large enough such that the hook parts overlap each other in phase space. The oxygen phase space density also shows multiple peaks in front of the DF. These peaks contain oxygen from different sources, including the reflected pre-existing current sheet oxygen, the lobe oxygen heated by the Hall electric field $E_z$, and the heated lobe oxygen from the exhaust. By tracking the oxygen trajectories, we found that the DF pushes the pre-existing current sheet oxygen ions by the Lorentz force, and it pushes the lobe oxygen ions in the exhaust by the electric field $E_x$ in the DF layer.

In addition, by comparing three different runs with 50%, 5%, and zero oxygen concentrations, we found that DFs formed with different thicknesses in the three runs.
Considering the pressure balance, we speculate that the DF thickness is mainly determined by the penetration depths of the pre-existing current sheet ions in the DF. Given that the major force on the penetrating ions by the DF is the Lorentz force, we estimated the penetration depth by considering the effective time that the force needs to reflect the ions away from the DF layer. The calculated penetration depths are close to the DF thicknesses in all three runs. This result demonstrates how the scale of DFs is determined by the composition of the pre-existing current sheet ions.

With the capability of observing oxygen velocity distribution functions at a high time resolution in a low-inclination orbit, the Magnetospheric Multiscale (MMS) mission is well-suited to observe oxygen dynamics related to magnetotail reconnection. The results of this chapter, including the acceleration of the lobe oxygen ions by the Hall electric field in the downstream outflow region, the hook-shaped reflected beam of oxygen in phase space, the multiple peaks in the oxygen distribution function in front of a DF, and the relation between the DF thickness and the ion composition in the current sheet, should be observable by MMS.
Chapter 6

Observation of Lobe O\(^+\) Acceleration

In Chapter 5, the O\(^+\) acceleration near and far away from the X-point was discussed based on implicit PIC simulations. One of the results shows that the lobe O\(^+\) ions located far away from both the X-point and the neutral sheet can be accelerated by the Hall electric field. In this chapter, we show a preliminary evidence of this effect by using spacecraft observations.

6.1. Spacecraft and Instruments

We studied a storm-time substorm event observed by Cluster spacecraft on 1 October 2001. This event has been discussed in several previous studies by different authors [e.g., Wygant et al., 2005; Kistler et al., 2005; Runov et al., 2006; Xiao et al., 2007; Eastwood et al., 2010]. In particular, we use this event to discuss lobe O\(^+\) acceleration far away from both the X-point and the neutral sheet. The data were obtained from the following instruments onboard Cluster: magnetic field from the Flux Gate Magnetometer (FGM) [Balogh et al., 2001], ion three dimensional distribution functions and moments with energy ranges from 40 eV/e to 40 keV/e from the time-of-flight ion Composition Distribution Function (CODIF) instrument in the Cluster Ion Spectrometry (CIS) package [Rème et al., 1997], and two dimensional vector electric field from the double probe Electric Field and Wave (EFW) instrument [Gustafson et al., 1997]. The GSM coordinate system is used in this chapter.
6.2. Results

6.2.1. Event Overview

Between 0920 and 1030 UT on 1 October 2001, the Cluster spacecraft was at X=−16 Re in the magnetotail plasma sheet. Figure 6-1 shows an overview of this event by Cluster C3 at [-15.9, 7.8, 1.1] Re. As indicated by the large Bx, C3 was far away from the current sheet in the intervals around 0920-0938 UT and 0955-1010 UT. There are multiple Bx reversals between 0938-0955 UT, indicating the current sheet crossings. In panel d, the H\(^+\) outflow \(V_{H_x}\) is mainly tailward before 0940 UT and mainly earthward flows after 0948 UT (there are a couple of flow reversals around 0948 of very short duration). The O\(^+\) outflow \(V_{O_x}\) is similar to \(V_{H_x}\) but with smaller flow magnitudes. The flow reversals indicate a crossing of the X-line. The X-line crossings are accompanied by small Bx indicating that the spacecraft is close to the reconnection sites, which are around 0940 UT and 0948 UT. In previous studies [e.g., Wygant et al., 2005; Eastwood et al., 2010], the ion diffusion region and ion acceleration have been studied. The electric field fluctuates strongly when the spacecraft is near the reconnection sites. Note that the EFW instrument detects only two dimensional electric fields in approximately the X and Y directions. The third component is calculated by assuming \(\vec{E} \cdot \vec{B} = 0\). Although this assumption might not be valid close to the reconnection sites, we will stay with this assumption since we are interested in the region far away from the reconnection sites in this chapter. The ion number densities in panel c show that the O+ species became significant after 0938 UT and dominant at around 0948-0955 UT. The O+ temperature is larger than the H+ temperature during the entire event.
Figure 6-1 Cluster C3 overview 0920-1030 UT on 1 October 2001. (a) Magnetic field components (b) electric field components (c) ion number density (d) H⁺ velocity components (e) O⁺ velocity components (f) ion temperatures.

6.2.2. Criteria for Selecting the Interval of Interest
Here, we discuss the criteria for selecting the interval of interest. The goal of this chapter is to investigate the lobe O\(^+\) acceleration by the Hall electric field near the separatrices far away from both the X-line and the current sheet by using the observations and compare the signatures with the simulation results in Chapter 5. The key point is how to determine that the observations are “far” from the observation X-line and current sheet. In the simulation in Chapter 5, we found that both the H\(^+\) and O\(^+\) are accelerated in the X direction in the exhaust. We also found that in the downstream outflow region, the H\(^+\) ions flow more quickly than the O\(^+\) due to the much larger inertia of the O\(^+\) so that in the exhaust ‘far” from the X-line only H\(^+\) ions are found flowing along the field. By contrast, the lobe O\(^+\) ions are accelerated locally in the Z direction by the Hall electric field near the separatrices and become a hot population in that region. Therefore, the first criterion is that \(B_x\) have large absolute values so that the location is far away from the current sheet and could be close to the separatrices. The second criterion is that the H\(^+\) ions flow along the magnetic field and not the O\(^+\). We can see this in Figure 6-2 which shows the simulated flow signatures for the \(V_{H,||}\) and \(V_{O,||}\) from the simulation in Chapter 5. Note that for both panels a and b, the Earth is on the right hand side. The green and blue circles near the separatrices indicate the locations in the outflow region “close” to and “far” away the X-point respectively. As we argued above when the location is “close” to the X-point, both the \(V_{H,||}\) and \(V_{O,||}\) are large. When the location is “far” away from the X-point, only \(V_{H,||}\) shows strong outflow while the \(V_{O,||}\) is much smaller. The third criterion is that there be a non-zero electric field in Y-Z plane as expected for the Hall electric field.
Figure 6-2 Ion velocity components parallel to the local magnetic field for (a) H⁺ and (b) O⁺. The green and blue circles indicate the locations “close” to and “far” away from the X-point. See the text for the details.

6.2.3. Flows and Velocity Distribution Functions
By using the criteria above, the region of interest is found to be between 10:05:20-10:05:40 UT is highlighted by the double vertical dashed lines in Figure 6-3. The velocities in the bottom three panels are in field-aligned coordinates: “||” is parallel to the local magnetic field; “⊥1” is the perpendicular direction in the plane determined by the H+ flow and magnetic field; “⊥2” follows the right hand rule. For “⊥1” and “⊥2”, the $E \times B$ velocity components are shown for comparison. In this interval, the $B_x$ is large and is the dominant magnetic field component, which indicates that the spacecraft is far away from the neutral sheet. In the last panel, $V_{H,||}$ is much stronger than $V_{O,||}$, which indicates this location is in the outflow region and far away from the X-line. The electric field in panel b shows a strong $E_z$ component essentially pointing toward the current, which is characteristic of the Hall electric field. In panel c, the $V_{H,⊥1}$ is consistent with the $E \times B$ velocity component in this interval, while both the $V_{O,⊥1}$ and $V_{O,⊥2}$ deviate from the $E \times B$ velocity components in both panels c and d. This is similar to the feature we discussed in Chapter 5 (see Figure 5-3 and 5-4 and related text) that the O+ ions are nonadiabatically accelerated by the Hall electric field, while the H+ ions are frozen-in and perform $E \times B$ drift in Y direction. Because the H+ ions drift in the Y direction, the “⊥1” is mainly in the Y direction and the “⊥2” is mainly in the Z direction. Thus, the acceleration by Hall electric field should be in “⊥2” direction. Note that the $V_H$ does not have a “⊥2” component, because the “⊥2” direction is perpendicular to the $V_H$ by definition. Also note that we cannot see clear bulk acceleration for O+ ions in “⊥2” direction in panel d. This is because there is a hot O+ population accompanying the beam with bulk acceleration in the velocity distribution. As a result, the average O+ flow in “⊥2” direction is small. These will be shown in Figure 6-4.
Figure 6-3 Cluster C3 observation for (a) and (b) magnetic and electric field components (c) and (d) “\(\perp 1\)” and “\(\perp 2\)” components of \(H^+, O^+\), \(E \times B\) velocity (e) “\(\parallel\)” component of \(H^+\) and \(O^+\) velocity. See the text for the definition of the field-aligned coordinates. The double dashed lines highlight the interval of interest.

Figure 6-4a and b shows the \(O^+\) velocity distribution function integrated in the highlighted interval in Figure 6-3. The distribution functions are shown in the \(\parallel\)-\(\perp 1\) and \(\perp 1\)-\(\perp 2\) slices. The red arrow, blue line, and pink line are the average electric field direction, \(O^+\) flow, and \(E \times B\) velocity projected in the slices respectively. The electric field is \(E(\perp 1, \perp 2, \parallel) = (1.11, 2.75, 0.01)\)
mV/m, where the “\(\perp 2\)” component is dominant as we expected. From both panels a and b, we can see there is a dense population approximately aligned with the electric field and some hot diffusive distributions around. The dense population shows the properties of the non-adiabatic acceleration of lobe cold O\(^+\) ions by Hall electric field. The hot population may be accelerated at an earlier time, at the same location or at another locations. The average O\(^+\) flow (i.e., the first moment) does not show any signatures of these two very different populations. Note that the observed distribution function plots (a) and (b) show a lot of “hollow” regions (white color). These are mainly caused by low particle counts. Figure 6-4c and d show the O\(^+\) distribution function slices in the simulation in Chapter 5 and they are plotted in the same format as panels a and b. Note that in the simulation, there is only one lobe O\(^+\) population with 5keV in the initial set up. Although this O\(^+\) population is hotter than the observed cold beam, they both show similar bulk acceleration by the electric field.
Figure 6-4 Slices of the O\textsuperscript+ velocity distribution functions in field-aligned coordinates (a) $V_\parallel$-$V_{\perp 1}$ and (b) $V_{\perp 1}$-$V_{\perp 2}$ observed by Cluster C3. They were obtained by integrating the distribution functions in the interval of interest, i.e., the interval in Figure 6-3 highlighted by the double dashed lines. The red arrows are the electric field projected on the slices. (c) and (d) are from the simulation in Chapter 5 and are plotted in the same format as (a) and (b). The blue line and the pink line in (a) and (b) are the O\textsuperscript+ flow velocity and the $E \times B$ velocity projected to the slices respectively.

### 6.3. Discussion and Summary
In this chapter, we show a preliminary observation of lobe O$^+$ acceleration by the Hall electric field far away from both the X-point and the neutral sheet. The Cluster event on 1 October 2001 is studied. Using the results from the simulation in Chapter 5, we designed some criteria to select the interval of interest from the event. By comparing the distribution functions from the observations with those from the simulation, we found that:

(1) Both the observation and simulation show similar bulk accelerated O$^+$ population aligned in the Hall electric field direction, which confirms that the lobe O$^+$ ions can be accelerated locally far away from the X-point as we discussed in Chapter 5.

(2) There is a diffusive hot O$^+$ population found in the observations which does not appear in the simulated distribution function.

Note that although the results by Cluster C3 are shown in this chapter, the C1 has been examined and shows similar results. For conclusion (1), although the observed beam is colder than the simulated one, this is due to the initial O$^+$ temperature set-up in the simulation and will not change the conclusion. For conclusion (2), the observed hot O$^+$ population could be the lobe O$^+$ ions accelerated earlier. It might also come from other reconnection sites or due to some instability in Y direction. These factors are not included in the 2D simulation in Chapter 5. A 3D simulation is needed in order to study such effects, which can be carried out in the future.

These are just preliminary observations. More detailed observations are need to quantify the effects of the lobe O$^+$ acceleration far away from the reconnection sites. With the capability of observing multi-ion species at multi-points, the improved instrumentation on Magnetospheric Multiscale (MMS) spacecraft should provide much better opportunities to observe such acceleration.
CHAPTER 7

Conclusions and Future Plans

In this dissertation, the electron and ion acceleration associated magnetotail reconnection in the Earth's environment is studied. Geomagnetic substorms are important for studying magnetotail energy release but vary significantly from event to event. Two factors affecting substorms are considered: (1) realistic magnetotail X-line configuration based on upstream solar wind conditions; (2) the presence of oxygen ions as a significant heavy ion species in the magnetotail. The associated acceleration of electrons and ions are discussed separately in this dissertation.

7.1. Conclusions on Electron Acceleration in Realistic Magnetotail X-line Configuration based on Upstream Solar Wind Conditions

The electron acceleration near the magnetotail X-line based on realistic magnetotail configurations driven by observed solar wind conditions is studied in Chapter 3. This study is carried out by using the UCLA global MHD model and LSK simulations. The MHD model simulates the response of the magnetotail under the observed solar wind conditions during specific substorm events. It provides the time-evolution of magnetotail dynamics including the electromagnetic fields, flows and the macroscopic configurations of the plasma sheet and the magnetotail X-line. The LSK simulation is used to calculate the electron behavior in the magnetotail electric and magnetic fields provided by the MHD model. In order to involve the adiabatic and non-adiabatic effects efficiently, it is designed to switch between the guiding
center and the full particle calculation schemes based on the value of the adiabaticity parameter \( \kappa \) which is defined as the square root of the ratio of the radius of curvature to the particle gyro-radius. This calculation is carried out for each particle at each time step. In this way, tens of thousands of particles are impulsively launched from near the magnetotail X-line and their motions are calculated individually. Once they hit virtual detectors placed throughout the simulation domain, information about their position, velocity, energy and the adiabaticity parameter are recorded. With such information, distribution functions can be calculated. In addition, individual particles can be selected and the acceleration processes acting on them can be evaluated. The fields and flows from the MHD model and the velocity distribution functions from the LSK simulation can be validated by comparing with DF events observed by spacecraft at specific locations and times during the substorms.

In this work, two substorm events, on 15 February 2008 and on 15 August 2001, were studied. Both the MHD and the LSK simulations of the two events were validated by comparing with observations from THEMIS and Cluster spacecraft. The main conclusions are as follows:

(1) According to the MHD model, the two events show very different magnetotail configurations, which are caused mainly by the different upstream solar wind conditions and the Earth's dipole tilt. The 15 February 2008 event had localized neutral lines accompanying a narrow (< 5 \( R_E \) in the cross-tail direction) earthward flow channel with relatively high speed (~ 300 km/s) flows in the magnetotail. By contrast, the 15 August 2001 event was characterized by cross-tail neutral lines and a broad region of (> 15 \( R_E \) in the cross-tail direction) earthward flows with relatively slow speed (~100 km/s).
Based on an LSK simulation, distribution functions at the locations earthward from the X-line for 15 February 2008, had $T_\perp / T_\parallel \approx 2.3$ at $X=-10$ to $-16$ RE, while for the 15 August 2001 event, the distribution functions had $T_\perp / T_\parallel = 0.8$ at $X=-10$ to $-15$ RE, where the $T_\perp$ and $T_\parallel$ are the temperatures in the parallel and perpendicular directions. The temperature was defined by using the second order moments of the distribution functions. These big differences were generally caused by variations in the combination of the adiabatic and non-adiabatic effects during their earthward convection due to the different magnetotail configurations for the two events.

The electron non-adiabatic acceleration and scattering can modify the distribution functions with temperature ratios from $T_\perp / T_\parallel = 1.0$ to $T_\perp / T_\parallel = 0.8$. These processes are influenced by the combination of the X-line configuration, the area of the regions with sharply curved field lines (could be small flux ropes) and the earthward flows. In this study 55.1% earthward propagating electrons experience non-adiabatic effects in the event with the cross-tail neutral line while only 0.5% of them are non-adiabatic in the event with localized neutral lines. The electrons experience non-adiabatic acceleration and scattering in the regions with sharply curved field lines, which can be qualitatively defined as the regions with low adiabaticity parameter $\kappa$ for the thermal electrons in the plasma sheet during the substorm ($\kappa<10$ for 2-keV electrons in this study). Fast earthward flows can convect electrons away from locations close to the X-line or with sharply curved field lines to earthward regions having strong magnetic field. Flows reduce the electron non-adiabatic effects by decreasing the time that they stay in the regions with low adiabaticity parameter $\kappa$. 

(4) The differences between the magnetotail configurations during these two substorms result from differences in the location of dayside reconnection. On 15 August 2001 the reconnection occurred on the flanks of the magnetopause, while on 15 February 2008 the reconnection occurred nearer the subsolar point.

7.2. Conclusions on Ion Acceleration by Magnetotail Reconnection in the presence of Oxygen Ions

Motivated by the typical observation of a significant fraction of oxygen ions near the magnetotail X-line during enhanced geomagnetic activity, the effects of oxygen ions on magnetotail reconnection and the ion acceleration in such reconnection are studied in Chapters 4 and 5. The study was carried out self-consistently by using a 2.5D implicit Particle-in-Cell simulation (iPIC3D). Compared to the general explicit PIC scheme, the stability constraints or the implicit scheme do not require us to resolve the temporal scale of the electron motion and spatial scales less than the Debye length. By properly selecting the time step and the grid spacing, the implicit PIC scheme can keep the self-consistent electron physics without the need for resolving electron scales. This method is an efficient way to study the long-time and large-scale effects (e.g., heavy ion scales) during magnetic reconnection in which the electron physics is important. To mimic magnetotail reconnection, a Harris equilibrium without a guide field was set up initially with a half-thickness equal to half of a proton inertial length. The same current sheet thicknesses are used for simulation runs with different oxygen concentrations. Both the theory of Harris current sheets [e.g. Galeev, 1982] and observations [e.g. Liu et al., 2014] indicate that the current sheet thickness is determined by the ambient magnetic field, the plasma temperature and the cross-tail speed [e.g., Galeev, 1982] and not the O+ concentration.
We carried out three simulation runs with different O\(^+\) concentrations: (1) “O\(^+\) Run”, with 50% O\(^+\) ions, i.e., the density ratios with respect to the electron density were \(n_e : n_H : n_O = 1 : 0.5 : 0.5\); (2) “EMD Run”, with equal mass density between protons and O\(^+\) ions, i.e., about 5.9% O\(^+\) ions; (3) “H\(^+\) Run”, with only electrons and protons. In the runs with O\(^+\) ions, the O\(^+\) ions occupy both the lobes and the current sheet initially as along with electrons and protons. The density in the lobe is uniform and is one tenth of the peak value of the current sheet density. The temperature ratio \(T_e : T_H : T_O = 0.2 : 1 : 1\) is based on the observations in the plasma sheet. The mass ratios of electrons, protons and oxygen ions with respect to proton are 1/128:1:16 respectively. The reconnection was initially triggered by a perturbation at the center of the simulation box which decayed exponentially with distance away from the center in order to avoid generating artificial box-scale waves. Open boundary conditions are used which allow particles to enter and leave the simulation box. The non-adiabatic acceleration mechanisms for the ions from the lobe and the pre-existing current sheet was studied by examining the phase space density, the velocity distribution functions, and by self-consistently tracing particles in the PIC simulation. The simulated lobe O\(^+\) acceleration is preliminarily confirmed by a substorm observation by Cluster on 1 October 2001 (Chapter 6).

The main conclusions in this study are:

(1) Based on the analysis of the force balance and plasma composition at the DF, if the current sheet includes a sufficient concentration (50% in this study) of O\(^+\) ions, the DF must overcome the large current sheet inertia due to the presence of the O\(^+\) ions. In the 50% O\(^+\) run, this effect reduces the DF speed by a factor of 2 and delays the fast reconnection phase by a factor of 2.
(2) The difference between the O\(^+\) and electron flow speeds in the O\(^+\) diffusion region generated an ambipolar electric field. The propagation speed of the magnetic flux is reduced by the drag of the O\(^+\) ions via this ambipolar electric field in the O\(^+\) diffusion region, which results in a smaller reconnection rate.

(3) DF thickness is related to the average penetration depth of the pre-existing current sheet ions. It is proportional to the O\(^+\) concentration in the pre-existing current sheet. The thickness is in the range between the local ion inertial lengths of proton and O\(^+\). It depends on which pre-existing ion species has the major contribution to the average penetration depth.

(4) The acceleration of ions from the lobes is mainly due to the electric fields in the diffusion regions of the reconnection sites. The electric fields in the exhaust, i.e., the outflow region between the X-point and the DFs, include the Hall electric field near the separatrices (\(E_x\) and \(E_z\) components), the reconnection electric field (\(E_y\) component), and the ambipolar electric field (\(E_x\) in the outflow region and \(E_z\) in the inflow region) in the O\(^+\) diffusion region. Here, the coordinates are in GSM coordinates. In the downstream outflow region, i.e., the outflow region downstream of a propagating DF, the electric field is mainly the Hall electric field (\(E_z\) component) near the separatrices.

a) For the lobe protons, 67.3% of them enter the exhaust and gain energy mainly from the reconnection electric field and the \(E_x\) component of the Hall electric field. The rest 32.7% of them enter the downstream outflow region and undergoes \(E \times B\) drift along the Y-direction near the separatrices without gaining energy. In the exhaust, the protons form a counter-streaming feature in the velocity distribution functions due to the \(E_z\) component of the Hall electric field. Such distribution
functions are hotter than the initial Maxwellian distribution function. The heated proton population in the downstream outflow region consists of the heated protons in the exhaust which flow downstream along the field lines.

b) For the lobe O\(^+\) ions, they enter the exhaust and gain energy from the electric fields as the protons do with the exception that they have a much larger diffusion region for the non-adiabatic acceleration. In addition, unlike the protons, the rest (42.4\%) of the lobe O\(^+\) ions directly enter downstream of the outflow region, far away from both the X-point and the neutral sheet and are accelerated by the Hall electric field near the separatrices. Because of this acceleration, they are able to move around the center of the neutral sheet and become a heated-oxygen component that populates the downstream outflow region in front of the DFs without entering the exhaust. In sharp contrast to the protons, because of their large inertia, few oxygen ions that are heated in the exhaust are found in the downstream outflow region.

(5) The pre-existing current sheet ions are reflected by the propagating DFs due to the Lorentz force. The dependence of the phase space density on \(V_x\) and \(X\) shows very different reflection signatures for the O\(^+\) ions and protons. For the protons, there is a reflected beam in front of the DF which has been observed by spacecraft [Zhou et al., 2010]. The O\(^+\) ions show a reflected beam with a hook shape in front of the DF. This hook-shaped reflected beam results from the increasing speed of the DF and can be used to deduce the DF's speed history during the simulation. This hook signature is difficult to identify for protons, because the typical proton thermal velocity in the tail plasma sheet is comparable to the reflection speed so that if there are hook like parts,
they will overlap each other in phase space. The oxygen phase space density also shows multiple peaks in front of the DF. These peaks contain oxygen from different sources, including the reflected pre-existing current sheet oxygen, the lobe oxygen accelerated by the Hall electric field $E_z$, and the heated lobe oxygen from the exhaust. (6) As found in the substorm event observed by Cluster on 1 October 2001, the $O^+$ velocity distribution functions, far away from both the X-line and the current sheet, have the bulk acceleration signature consistent with acceleration by the Hall electric field. This confirms the previous conclusion (4) (b) on the lobe $O^+$ acceleration in the simulation. Besides the bulk-accelerated population, there are some diffusive hot $O^+$ measured by spacecraft that are not generated the simulation. This hot population might come from other reconnection sites or due to some instabilities in Y direction that are not included.

7.3. Future Plans

The magnetotail reconnection in the presence of the $O^+$ ions and the associated ion acceleration were studied by using implicit PIC simulations in Chapters 4 and 5. The next step in this study is to test the simulation results by comparisons to observations and apply the results to improve understanding of the observations. To this end, the future plan is:

(1) The simulation in the presence of the $O^+$ ions shows several properties of DFs including the slow DF propagation, the thick DF thickness, the ambipolar electric field in the $O^+$ diffusion region, the $O^+$ acceleration by the Hall electric field far away from the X-point and the neutral sheet, the heated $O^+$ population in front of the DF, and the hook-shaped reflected beam formed by the propagating DF in $O^+$ phase space. All these properties are needed to compare with the spacecraft observations. Because the
time resolution for the O\textsuperscript{+} ions is low (~8-16s) for recent missions, i.e., Cluster and Magnetospheric Multiscale (MMS) spacecraft, the observed O\textsuperscript{+} reflected beam on one spacecraft might have only one data point. Therefore, multi-point observations by Cluster and MMS will be used to resolve the O\textsuperscript{+} reflected beam.

(2) The signatures of non-adiabatic acceleration of O\textsuperscript{+} in the velocity distribution functions by the electric field associated with the reconnection can be used to detect reconnection in spacecraft observations. Due to the large gyro-radii of the O\textsuperscript{+} ions, the region with O\textsuperscript{+} non-adiabatic acceleration is large and thus the reconnection signal can be remotely sensed. As shown in Chapter 5, O\textsuperscript{+} ions heated by the Hall electric field can be found in the quiet current sheet in front of the DFs. In order to apply this idea to remote-sensing of the reconnection, more simulation studies (in 2D, 3D and realistic magnetotail configuration) need to be carried out to build up a database of the O\textsuperscript{+} distribution functions at different locations related to the reconnection site. This database can be used to compare and explain the observations.

(3) The evolution of the reconnection can be deduced and studied from the O\textsuperscript{+} phase space plots. Because the O\textsuperscript{+} gyro-period is comparable to the evolution time scale of the reconnection, the signatures in O\textsuperscript{+} phase-space caused by the evolving electromagnetic field around the reconnection site reflect the history of the reconnection. Examining such signatures in observations can provide time-varying information about the reconnection. This can improve our understanding of the evolving reconnection beyond the steady reconnection models. Similar to (2), more simulations are needed for better understanding of the connections between the O\textsuperscript{+} phase-space signatures and the reconnection processes. Also, it is necessary to investigate the upper limit of the
time of the evolution that can be deduced in this way and the range of the distances from the reconnection site over which this technique can be used.

(4) An ambipolar electric field is found in the O\textsuperscript{+} diffusion region. This ambipolar electric field added O\textsuperscript{+} drag to the propagating magnetic flux and thus reduced the reconnection rate. This effect needs to be determined quantitatively in order to find out its actual contribution to the decrease in the reconnection rate. In addition, how this effect depends on the concentration of the O\textsuperscript{+} ions is unclear.
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