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Magnetic Reconnection for Coronal Conditions: Reconnection Rates, Secondary Islands and Onset

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Abstract Magnetic reconnection may play an important role in heating the corona through a release of magnetic energy. An understanding of how reconnection proceeds can contribute to explaining the observed behavior. Here, recent theoretical work on magnetic reconnection for coronal conditions is reviewed. Topics include the rate that collisionless (Hall) reconnection proceeds, the conditions under which Hall reconnection begins, and the effect of secondary islands (plasmoids) both on the scaling and properties of collisional (Sweet-Parker) reconnection and on the onset of Hall reconnection. Applications to magnetic energy storage and release in the corona are discussed.

Keywords Magnetic reconnection · Collisionless reconnection · Hall reconnection · Secondary Islands · Plasmoids · Solar corona

1 Introduction

A long-standing question about the solar corona is why it is nearly two orders of magnitude hotter than the solar surface. (See, *e.g.*, Klimchuk 2006 for a review.) Understanding coronal heating is important for many reasons, including its role in the acceleration of the solar wind. Competing models of coronal heating include wave heating and heating by nanoflares. In the nanoflare model, small flares releasing as little as a billionth as much energy as large flares occur essentially continuously both in active regions and the quiet corona (Parker 1983, 1988). The eruptive release of energy in flares is likely caused by a change of magnetic topology during magnetic reconnection (Priest and Forbes 2002). Studying the fundamental physics of magnetic reconnection can contribute to our understanding of these coronal processes.

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In addition to coronal heating, reconnection is important in other settings. Recent review articles have addressed many examples: large flares and coronal mass ejections (Priest and Forbes 2002; Isobe and Shibata 2009), magnetospheric convection and geomagnetic storm and substorm phenomena for space weather applications (Eastwood 2008), particle acceleration (Mozer and Pritchett 2010), and astrophysical, fusion, and laboratory settings (Zweibel and Yamada 2009).

The present review article focuses mainly on three topics: (1) Can magnetic reconnection release energy as fast as required by observations of flares? (2) Why does magnetic reconnection in flares begin abruptly? In other words, how does reconnection play a role in the storage and release of free magnetic energy? (3) How do recent developments on secondary islands in collisional (Sweet-Parker) reconnection affect our understanding of reconnection onset for coronal conditions?

This review article is organized as follows. A brief discussion of some observational constraints and the way in which numerical simulations can be used to study reconnection is in Sect. 2. Classical models of magnetic reconnection are summarized in Sect. 3. Recent results about reconnection onset in weakly collisional systems such as the solar corona are discussed in Sect. 4. Section 5 summarizes recent results on secondary islands in Sweet-Parker reconnection and its impact on reconnection onset. Section 6 contains a discussion of how the recent developments summarized here impact our understanding of coronal energy storage and release, as well as a discussion of some assumptions that go into the models and remaining open questions.

2 Observational Constraints and Reconnection Simulation Methodology

Many energy release events believed to be produced through reconnection exhibit two disparate timescales: a long period during which magnetic energy builds up but very little energy is released which we refer to as the "build-up phase" and a period of significant energy release which often begins suddenly which we call the "energy release phase." The key point is that most of the magnetic energy is released in the latter stage. Examining the build-up and energy release phases in the context of two space plasma systems yields interesting insights. In typical X-class solar flares, a large energy release in the form of X-rays and electron energization occurs for a period of about 100 seconds (Miller et al. 1997), but an active region on the sun may exist for weeks without producing such a flare. During a magnetospheric substorm, a significant fraction of lobe flux is reconnected in a period of about 10 minutes causing a dipolarization of the magnetotail and particle energization that creates the aurora, but the typical time between repeating substorms is about 3 hours (Borovsky et al. 1993).

There have been many ideas put forth on why these eruptions onset abruptly. In this review, we focus on coronal applications rather than magnetospheric applications. While there are many similarities between the two (Reeves et al. 2008; Lin et al. 2008; Linton and Moldwin 2009), they are different in that the corona is ostensibly weakly collisional before onset while the magnetosphere is essentially collisionless, which likely implies a different onset mechanism in the two settings.

Models of the onset of flares and coronal mass ejections (CMEs) can be divided into two sets: ideal-magnetohydrodynamic (MHD) driven and reconnection driven. Examples of ideal-MHD driving are the kink instability (Hood and Priest 1979), a catastrophic loss of equilibrium (Forbes and Isenberg 1991), and the so-called breakout model (Antiochos et al. 1999). Examples of reconnection driven models include tether cutting (Sturrock 1989), the tearing instability (Kliem 1995), and the onset of fast magnetic reconnection due to an abrupt beginning of Hall reconnection (Ma and Bhattacharjee 1996; Bhattacharjee 2004; Cassak et al. 2005) or other kinetic micro-instabilities (Sato and Hayashi 1979). It is important to note that the large number of models do not imply that one is right and the others are wrong; observations have revealed evidence supporting different models in different flares and CMEs (Sterling and Moore 2004). It appears that the question is not "Which is the right model?", rather "Which is the right model for a particular circumstance?"

If reconnection or tearing related phenomena is the driver of the onset of eruptions, it is obvious that the microphysics of reconnection is important for such applications. However, even if ideal-MHD effects are driving the onset of the eruptions, understanding the microphysics of reconnection and its onset is important. The examples of breakout and the catastrophic loss of equilibrium rely on reconnection to complete the process, so understanding how it occurs in such settings is necessary. In particular, a thin current sheet is formed before the eruption in both models and must be persistent for an extended period of time before an eruption occurs. If studying the microphysics of reconnection reveals that such current sheets are not stable or reconnect rapidly, then such models may need further scrutiny. For the kink instability, some have argued that reconnection is not necessary to allow an eruption (Török and Kliem 2005). However, some flux ropes undergo a kinking but do not erupt (so-called "failed eruptions"). It seems likely that reconnection, along with other effects such as the overlying magnetic field, plays a role in determining whether an eruption occurs. For these reasons, understanding the microphysics of reconnection may help provide insight into models of eruptions that take place on the larger scale. The present review focuses on recent developments on this topic and how it impacts models of solar activity.

The difficulty with understanding how reconnection occurs in coronal settings is the extremely disparate length and time scales involved. In collisionless systems, the diffusion region has length scales associated with the electron and ion skin depths, c/ω_{pe} and c/ω_{pi}. With a rough estimate of the density in the solar corona ($n \approx 10^{10}$ cm^{-3}), these scales are $c/\omega_{pi} \sim 200$ cm and $c/\omega_{pe} \sim 5$ cm. A typical magnetic flux tube involved in a flare has a volume of 10^{27} cm^3. (It has long been assumed that collisionless effects are not important in the corona because of the small length scales, but we argue in the following section that they are not negligible for reconnection.) Performing numerical simulations that resolve both diffusion region length scales and system size scales simultaneously is far beyond the capability of computers for the foreseeable future.

How can progress in numerical studies be made in understanding whether and how magnetic reconnection is the physical mechanism releasing this energy? One characteristic that helps is that the build-up of magnetic energy in these systems is believed to be due to external flows which compress the magnetic flux and create and intensify current sheets (Parker 1983, 1988). Eventually, these current sheets become unstable to magnetic reconnection, releasing the stored energy. A key point, however, is that the disparate timescales of the build-up and release phases are evidence that the external forcing flows are much slower than the inflows due to magnetic reconnection, *i.e.,* $v_{force} \ll v_{in}$. Just before magnetic energy release, the system can be justifiably pictured as a static equilibrium with a long thin current sheet. In order to simulate these systems realistically, it is not necessary to provide external forcing to the system; simulating a current sheet unstable to reconnection release is sufficient.

Initial value simulations of such long thin current sheets in collisionless systems typically also reveal two distinct phases: a long time over which a finite magnetic island slowly forms and flows develop which we call the "developmental phase," and a fast reconnection phase which we call the "asymptotic phase." The key is that almost all of the energy release occurs when reconnection has reached its asymptotic phase (Shay et al. 2004), so one can focus on the physical scaling of the energy release during the asymptotic phase. Because of the small size of the island and very weak flows during the developmental phase, in any real system the developmental phase will be indistinguishable from the build-up phase.

Gaining a physical understanding of the nature of the build-up and developmental phases is a very difficult, but rich, problem. The theory must predict rates fast enough to allow reconnection to grow from noise to significant size during the developmental phase, but it also must be slow enough to allow significant magnetic loading during the build-up phase between energy release events. Obtaining a theory that satisfies these timescale demands is difficult because the developmental phase is strongly dependent on the specific initial conditions and parameters in the system, and these conditions are to a large extent unknown, *e.g.,* the size of the perturbations or forcing creating the X-line (Wang et al. 2000), the thickness of the initial current sheet (White 1986), the size of the system (Grasso et al. 1999), the presence of a magnetic field normal to the current sheet (Lembege and Pellat 1982; Pellat et al. 1991), and the specific kinetic equilibrium of the electrons (Sitnov et al. 2002). The growth of reconnection from noise has also been shown to be strongly dependent on the electron to ion mass ratio (Porcelli 1991) and the resistivity (Furth et al. 1963). Note that the literature on the developmental phase of reconnection is extensive, and the citations above are only included to be illustrative rather than exhaustive. Citations to papers addressing other issues such as the dependence on kinetic effects are omitted. Compared to the developmental phase, the asymptotic phase of reconnection is much simpler, and we seek primarily to address the asymptotic phase in this article.

A fundamental question about the asymptotic phase is: is reconnection fast enough to explain the energy release time scales seen in physical systems? In the case of solar flares and substorms this question reduces to: is it possible for a significant fraction of the available magnetic flux to reconnect in 100 seconds and 10 minutes, respectively? In the context of this question, it is unimportant how long it takes the reconnection to initiate. If the conditions upstream of the diffusion region are changing much slower than the plasma transit time through the diffusion region, then a quasi-steady (Sweet-Parker-type) theory can be used to understand the reconnection.

The rate at which reconnection proceeds is quantified by the reconnection rate E, which is the out-of-plane electric field in the upstream and boundary layers and is usually normalized to $c_A B/c$, where c_A is the Alfvén speed based on the reconnecting magnetic field B immediately upstream of the diffusion region. From Faraday's law, E is a measure of the temporal rate of change of magnetic flux across the boundary layer. The reconnection rate during eruptive flares has been estimated from direct observations to be $E \sim 0.001$–0.2 (Ohyama and Shibata 1998; Yokoyama et al. 2001; Isobe et al. 2002; Qiu et al. 2002; Fletcher et al. 2004; Lin et al. 2005; Isobe et al. 2005). Therefore, reconnection can be a candidate for the energy release mechanism if E is in this range.

During reconnection, a non-magnetohydrodynamic (non-MHD) region called the ion dissipation region forms near the X-line with a half-length of L along the outflow direction and a half-thickness of δ along the inflow direction. The geometry of this region is very important for determining the scaling of reconnection rate; one can show that $E \propto \delta/L$. Thus, the key is to understand the physics controlling the aspect ratio of the diffusion region δ/L. If L scales like the system size $L_{\rm sys}$, the reconnection will be much too slow to explain the energy release during solar flares and substorms. If $L \ll L_{\rm sys}$, it is much more likely that reconnection can explain the release rates seen in nature. Therefore, one plausible approach to study whether reconnection is fast enough to explain observations is to (1) simulate a long thin current sheet which is strongly unstable to reconnection, and (2) examine the asymptotic behavior of the reconnection and determine the physics governing the aspect ratio of the diffusion region.

Even with this approach, there is significant uncertainty about the applicability of the conclusions of simulations to solar reconnection. The range of length scales is so large that it is possible that reconnection could be modified by new physical mechanisms operative at length scales much larger than system sizes attainable in current simulation studies. In addition, many simplifying assumptions are commonly employed in reconnection theories and simulation studies, such as

– two-dimensionality, with no variation perpendicular to the reconnection plane, and
– symmetry about the diffusion region in the up-down and left-right directions.

These assumptions cannot be expected to be valid in many naturally occurring settings, and work has been done to relax these assumptions (see *e.g.*, Priest and Schrijver 1999; Simakov et al. 2006; Mozer and Pritchett 2010, respectively). In particular, the assumption of two-dimensionality leaves out a wide range of instabilities which can disrupt and modify the reconnection current sheet. In the present discussion, all of these assumptions are employed, and relaxing them is discussed in Sect. 6.

3 Classical Models of Magnetic Reconnection

3.1 Sweet-Parker (Collisional) Reconnection

The Sweet-Parker model (Sweet 1958; Parker 1957) uses a steady-state scaling analysis to determine how reconnection parameters scale with system plasma parameters, which are assumed known. Sweet-Parker-type analyses have been used extensively to understand the reconnection process (Biskamp 1986; Wang et al. 1996; Uzdensky et al. 1996; Ji et al. 1998; Dorelli and Birn 2003). The Sweet-Parker model has been reviewed elsewhere (*e.g.*, Zweibel and Yamada 2009), so we do not derive it here. The result is

$$E \sim \frac{v_{\rm in}}{v_{\rm out}} \sim \frac{\delta}{L} \sim \sqrt{\frac{\eta c^2}{4\pi c_A L}}, \tag{1}$$

where $v_{\rm in}$ and $v_{\rm out}$ are the inflow speed into and outflow speed out of the collisional boundary layer, η is the resistivity, and c_A is the Alfvén speed based on the magnetic field B and plasma density n immediately upstream of the boundary layer. The term under the square root is the inverse of a Lundquist number S, though care should be used since L is the half-length of the layer, not the system size, and the Alfvén speed is based on the plasma properties immediately upstream of the reconnection layer, not necessarily characteristic global values.

The Sweet-Parker model was a major breakthrough because it is much faster than straight diffusion (Parker 1957). The model has been verified in laboratory experiments (Ji et al. 1998; Trintchouk et al. 2003; Furno et al. 2005) and numerical simulations (Biskamp 1986; Uzdensky and Kulsrud 2000). However, as has now been established, the model is only valid for Lundquist numbers below $S_{\rm crit} \sim 10^4$. There are drawbacks to the Sweet-Parker model which arise in many physical settings, even within the confines of the assumptions listed previously:

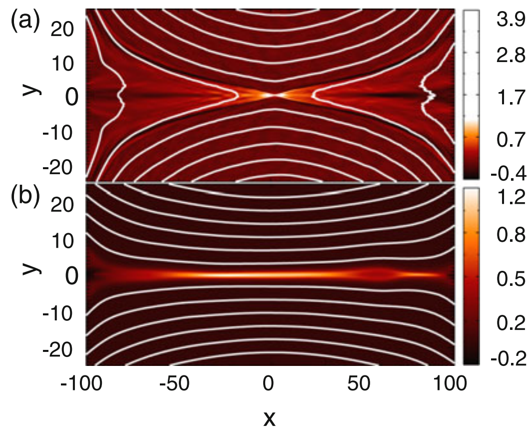

Fig. 1 Out-of-plane current density (*color plot*) with magnetic field lines overlaid (*white lines*) for (**a**) Hall reconnection and (**b**) Sweet-Parker reconnection, from (Cassak et al. 2010)

- For systems with large Lundquist numbers, (1) implies that E is small. This implies the process too slow to explain observed energy release times in flares, for which $S \sim 10^{12}$–10^{14}, and many other reconnection sites.
- For large S, (1) also implies that $\delta \ll L$, giving rise to elongated layers, such as seen in the out-of-plane current J_z plotted in Fig. 1(b). Highly elongated current sheets are unstable to a secondary tearing instability which gives rise to secondary islands (Biskamp 1986), also known as plasmoids upon ejection from the boundary layer. These alter the Sweet-Parker scaling results and make the reconnection faster (Matthaeus and Lamkin 1985; Lapenta 2008). This is the topic of Sect. 5.
- The Sweet-Parker model is based on resistive-MHD. Thus, it does not apply when length and time scales are small enough to give rise to other physics, such as the Hall effect and other kinetic effects (Vasyliunas 1975). Thus, the Sweet-Parker model is unlikely to explain any reconnection events in the Earth's magnetosphere or solar wind where the ion gyroradius exceeds collisional scales.
- Observed reconnection rates in the corona (close to 0.1) imply electric fields greatly exceeding the Dreicer runaway electric field (Drake and Shay 2007; Daughton et al. 2009b; Roytershteyn et al. 2010), at which Coulomb collisions between electrons and ions shut off (Dreicer 1959). Thus, it is essentially impossible for any model based on resistive-MHD with a Spitzer resistivity to explain the most rapid reconnection events in the corona. Collisionless physics beyond resistive-MHD, such as the Hall effect, is required.

It should be noted, however, that it has been proposed that Sweet-Parker reconnection can occur in the chromosphere (Litvinenko 1999; Chae et al. 2003; Litvinenko and Chae 2009), where Lundquist numbers are lower and kinetic physics may appear at length scales below collisional scales. Also, though collisional reconnection cannot be occurring during the energy release phase of a flare, it is not a priori precluded from occurring during the build-up phase (Shibata and Tanuma 2001; Cassak et al. 2005; Uzdensky 2007; Cassak et al. 2008). These topics are discussed further in Sect. 6.

3.2 Collisionless Reconnection

The consideration of reconnection beyond resistive-MHD was explored in a review paper by (Vasyliunas 1975). The linear tearing mode with kinetic effects was studied for fusion applications (Drake and Lee 1977; Terasawa 1983; Hassam 1984); it was determined that they increase the growth rate of the tearing mode. It was confirmed numerically (Aydemir 1992; Kleva et al. 1995) that reconnection is much faster when electron pressure is present. While some of these models contain collisions, Kleva et al. (1995) showed that the reconnection rate is fast when the thickness δ of the layer is below kinetic scales and collisions are playing no role. Many further studies found a similar enhanced rate of reconnection in the nonlinear regime (Wang and Bhattacharjee 1993; Zakharov et al. 1993; Rogers and Zakharov 1996; Ma and Bhattacharjee 1996; Cafaro et al. 1998; Grasso et al. 1999).

Since that time, many simulations have revealed that asymptotic reconnection rates in collisionless regimes approach a value of reconnection roughly independent of $(c/\omega_{\mathrm{pi}})/L$ and $(c/\omega_{\mathrm{pe}})/L$, *i.e.*, the reconnection rate is on the order of 0.1 and is independent of system size and electron mass (Shay and Drake 1998; Shay et al. 1999; Hesse et al. 1999; Birn et al. 2001; Huba and Rudakov 2004; Shay et al. 2007; Daughton 2010, and references therein). In this respect, bearing in mind that simulations can only resolve modest values of $(c/\omega_{\mathrm{pe}})/L$ and $(c/\omega_{\mathrm{pi}})/L$, there is general agreement that current simulation scaling implies that reconnection can be fast enough due to kinetic effects to have energy release times consistent with solar flares and magnetospheric substorms.

What is currently controversial is the physics responsible for these fast reconnection rates. It has been suggested (Mandt et al. 1994; Shay et al. 1999) that the Hall term is critical to allow these fast energy release rates. The Hall term operates at sub-ion gyroradius scales and describes the decoupling of ions from the magnetic field when their gyro-orbit is comparable to gradient scales in the magnetic field. Mandt et al. (1994) and Rogers et al. (2001) argued that collisionless (Hall) reconnection is fast because of the dispersive nature of the whistler and kinetic Alfvén waves introduced by the Hall term. This is because the increase in flow speeds at smaller length scales allows Petschek-type open outflow configurations (Drake and Shay 2007). Hall physics was found to fundamentally alter the properties of the diffusion region (Biskamp et al. 1995), and the quadrupolar magnetic field perturbations during kinetic reconnection (Sonnerup 1979) were shown to be a direct effect of electron physics due to the Hall term (Mandt et al. 1994). The importance of the Hall effect was highlighted in the GEM Challenge study (Birn et al. 2001 and references therein), which compared fluid, two-fluid, hybrid, and particle-in-cell (PIC) simulations. All simulations containing Hall physics had a similar (fast) reconnection rate, while the simulation without the Hall effect was much slower. Recently, large scale kinetic-PIC simulations have found asymptotic reconnection rates independent of electron mass and system size (Shay et al. 2007; Daughton 2010) as in previous studies. In addition, it was found that the structure of the outflow jet critical to fast reconnection rates is controlled by the Hall fields in the dissipation region (Drake et al. 2008), consistent with the Hall model.

There is much observational support for the occurrence of Hall reconnection. Signatures of Hall reconnection have been observed in the Earth's magnetosphere (Nagai et al. 2001; Øieroset et al. 2001; Scudder et al. 2002; Mozer et al. 2002; Runov et al. 2003; Borg et al. 2005; Phan et al. 2007) as well as those of other planets (Eastwood et al. 2008). It has also been studied in laboratory reconnection experiments (Ren et al. 2005; Cothran et al. 2005; Yamada et al. 2006; Frank et al. 2006). There is no direct evidence of Hall reconnection in the corona because the length scales involved (on the scale of meters) is far below the resolution of satellite observations.

However, the role of the Hall effect in allowing fast reconnection is not universally accepted. Karimabadi et al. (2004) performed hybrid simulations in which the Hall term was manually removed from the generalized Ohm's law, with resultant reconnection rates comparable to the case with the Hall term present being reported. A subsequent hybrid study, however, found the opposite result: removing the Hall term led to a long thin diffusion region with a slower reconnection rate (Malakit et al. 2009).

In large scale kinetic-PIC simulations (Karimabadi et al. 2005; Daughton et al. 2006; Karimabadi et al. 2007; Daughton et al. 2011), it was concluded that the length of the electron current layer tends to increase in time and is limited by secondary island formation or boundary conditions. The Hall effect alone would therefore not be sufficient to give fast reconnection. On the other hand, kinetic-PIC simulations also see significant long periods of relatively steady reconnection rates and diffusion region lengths (Karimabadi et al. 2007; Shay et al. 2007; Klimas et al. 2008). Further study is required to resolve the role of secondary islands in collisionless reconnection.

Beyond traditional plasmas (consisting of electrons and protons), simulations of reconnection in electron-positron plasmas can in principle provide a way to test the importance of the Hall effect, as in these systems the Hall effect identically cancels (Bessho and Bhattacharjee 2005). Simulations show that electron-positron reconnection is also fast (Bessho and Bhattacharjee 2005, 2007; Hesse and Zenitani 2007; Daughton and Karimabadi 2007; Swisdak et al. 2008; Zenitani and Hesse 2008), and therefore fast reconnection can occur in the absence of Hall physics. Any controversy arises from the general conclusions of these simulations relative to electron-proton plasmas. If a universal mechanism is allowing fast reconnection in both electron-positron and electron-proton plasmas, then Hall physics may not be playing the critical role even in the electron-proton case. Such universal processes could be secondary island formation (Daughton and Karimabadi 2007) or off-diagonal pressure tensor effects (Bessho and Bhattacharjee 2005; Hesse et al. 2009). On the other hand, it is possible that a different mechanism not present in the electron-proton case is allowing fast reconnection, in which case the cause of the fast rates in electron-positron reconnection would have no bearing on reconnection in traditional plasmas. One possibility discussed in the literature is the Weibel instability (Zenitani and Hesse 2008; Swisdak et al. 2008), which becomes active in the outflow jet and broadens out the current layer, though this remains a topic of debate (Daughton et al. 2011).

4 The Onset of Fast Reconnection in Weakly Collisional Systems

An important topic for understanding energy storage and release in the corona is what causes the onset of Hall reconnection in weakly collisional systems, where collisional reconnection may be important before onset. This ostensibly excludes reconnection onset in the magnetosphere where collisions are not expected to play any role, but should be relevant for coronal applications and possibly fusion devices.

One can estimate the length scale where MHD breaks down by comparing ideal-MHD terms to other terms in the generalized Ohm's law. If the field lines are anti-parallel, the characteristic length scale is c/ω_{pi}, the ion inertial length (equivalent to the Larmor radius for a particle traveling at the Alfvén speed) (Vasyliunas 1975). When there is a large out-of-plane (guide) magnetic field, the length scale associated with the Hall term becomes ρ_s, the Larmor radius for ions moving at the sound speed $c_s^2 = (T_e + T_i)/m_i$, where m_i is the ion mass and T_e and T_i are the electron and ion temperatures (Rogers and Zakharov 1995; Kleva et al. 1995). The question becomes—how does reconnection transition from Sweet-Parker reconnection at δ larger than kinetic scales to Hall reconnection at kinetic scales?

It was recently established that the transition from Sweet-Parker to Hall reconnection is abrupt (catastrophic) as a function of δ (Cassak et al. 2005). This occurs because both the Hall effect and the resistive term scale like $1/\delta^2$ (Birn et al. 2001), so that one or the other dominates for a given δ. It follows that magnetic reconnection in weakly collisional plasmas is bistable, as either term can dominate depending on the plasma parameters. This is true up to extreme values of the Lundquist number. The existence of abrupt transitions and bistability has been confirmed in numerical simulations both without (Cassak et al. 2005) and with (Cassak et al. 2007) a guide field. There is evidence from laboratory experiments that an abrupt transition occurs at $\delta \sim c/\omega_{\rm pi}$ (Ren et al. 2005; Yamada et al. 2006) or ρ_s (Egedal et al. 2007; Katz et al. 2010).

There have been recent attempts to perform a scaling analysis which describes the nonlinear dynamics of Hall reconnection (Chacón et al. 2007; Simakov and Chacón 2008; Malyshkin 2008; Tsiklauri 2008; Simakov and Chacón 2009; Zocco et al. 2009; Malyshkin 2009). While much has been learned, these models do not predict the length of the layer L (Uzdensky 2009; Sullivan et al. 2009), leaving current analytical models incomplete.

It has been suggested that the catastrophic onset of fast reconnection has important implications for why many reconnection events begin abruptly (Cassak et al. 2005, 2006; Uzdensky 2007; Cassak et al. 2008). The idea is that reconnection before the eruptive event is collisional, and the dynamics causes the layer to thin to kinetic scales, where reconnection onsets abruptly. While this model is appealing, it assumes that Sweet-Parker reconnection is valid for coronal conditions. However, as discussed in the previous section, the appearance of secondary islands alter Sweet-Parker reconnection at high Lundquist number. We now summarize recent developments about secondary islands in Sweet-Parker reconnection with an eye toward its impact on onset.

5 Effect of Secondary Islands on Collisional (Sweet-Parker) Reconnection

As discussed in Sect. 3, very elongated layers are predicted to occur during Sweet-Parker reconnection with high Lundquist numbers. When the sheet becomes sufficiently elongated, it breaks up due to a secondary tearing instability (Biskamp 1986). An example of a secondary island can be seen in Fig. 1(b) near the right side of the layer. As secondary islands and plasmoids appear in many different contexts, it is important to carefully delineate which manifestation of secondary islands is being treated in the present context. Here, we discuss secondary islands that form during collisional (Sweet-Parker) reconnection as a result of self-consistent evolution through a secondary tearing-type process. In particular, we are not treating secondary islands in the following contexts:

– As emphasized by Daughton et al. (2006), secondary islands can occur during collisionless reconnection. Islands have been observed during collisionless reconnection in the magnetosphere (Chen et al. 2008). The treatment here is before any transition to collisionless reconnection has occurred.
– The Sweet-Parker model tacitly assumes that the incoming plasma is laminar. If it is turbulent with sizable fluctuations, this could induce topology changes in multiple locations along the layer, resulting in plasmoids. This has been emphasized by Matthaeus and Lamkin (1985), Lazarian and Vishniac (1999) and studied in various contexts (Smith et al. 2004; Loureiro et al. 2009; Kowal et al. 2009; Nakamura et al. 2010) (although often the turbulence is imposed near the layer as opposed to the upstream region). This type of plasmoid formation is potentially very important, and will be discussed further in Sect. 6. In the present context, we will retain the simplifying assumption that the upstream flow is initially laminar.
– The secondary tearing instability appears to be a linear instability, and this phase has received recent attention. Linearizing around a dynamic equilibrium allowed a prediction (Loureiro et al. 2007) of the growth rate and number of islands. Predictions of this model

have been confirmed by numerical simulations designed to test the linear theory (Samtaney et al. 2009; Ni et al. 2010; Huang and Bhattacharjee 2010). It is to be suspected that the linear phase of the instability is short lived, and simulations have shown that the nonlinear phase differs in many ways from the linear phase (Daughton et al. 2009a; Cassak et al. 2009; Huang and Bhattacharjee 2010). The present treatment considers only the nonlinear phase of the instability.
– Finally, the assumptions discussed in Sect. 3 (two-dimensions and symmetric) are retained. Also, secondary island generation is manifestly time dependent, but we consider steady-state properties of the reconnection. By this, we mean that there is a continual generation and expulsion of plasmoids. When variations caused by secondary islands are averaged over sufficiently long times, quasi-steady properties emerge.

The recent work we will summarize addresses three main topics: (1) Does the resistive-MHD property of secondary island formation persist in more realistic descriptions of a reconnecting plasma, such as those described by PIC simulations? (2) What is the quantitative effect of secondary islands on Sweet-Parker reconnection? (3) Does the presence of secondary islands alter the way in which Hall reconnection begins and which mode dominates when both are simultaneously present?

5.1 Persistence of Secondary Islands in Kinetic Plasmas

The resistive-MHD description usually encountered in numerical simulations achieves closure through a resistivity η treated as a constant or as a Spitzer resistivity with a $T^{-3/2}$ dependence. However, they are not calculated self-consistently, so it is important to ensure that this description is applicable when applied to more accurate descriptions of a plasma, such as PIC codes. To address this, Daughton et al. (2009b) employed a Fokker-Planck collision operator into a PIC code, so that resistivity is self-consistently determined through the dynamics rather than being imposed. Also, Ohmic heating is automatically included, as well as the Dreicer runaway condition which is absent from fluid models.

Surprisingly, the results were largely consistent with resistive-MHD simulations, at least for the parameter regimes in which agreement would be expected (Daughton et al. 2009a, 2009b). In particular, when the layer is thicker than kinetic scales, a Sweet-Parker layer forms. For high enough Lundquist numbers, the Sweet-Parker layer self-consistently develops secondary islands. The main difference between the resistive-MHD and PIC descriptions is that the effective resistivity from the PIC simulations self-consistently evolves in time because of the temperature change in the plasma, which can cause a thinning of the layer.

5.2 Scaling of Sweet-Parker Reconnection with Secondary Islands

When secondary islands occur, the simple Sweet-Parker scaling law in (1) is no longer expected to hold. A quantitative prediction of the mean properties of the reconnection was not developed until recently. It was proposed that if the fragmented current sheets that form when secondary islands occur are still collisional, they should be described by Sweet-Parker scaling (Daughton et al. 2009b). This is analogous to a similar assumption for reconnection in turbulent plasmas (Lazarian and Vishniac 1999).

This model allows one to make relatively simple quantitative predictions of how reconnection with secondary islands proceeds. If a Sweet-Parker diffusion region of half-length L_{SP} has N secondary islands, it is effectively cut into pieces of average length $L \sim L_{SP}/N$.

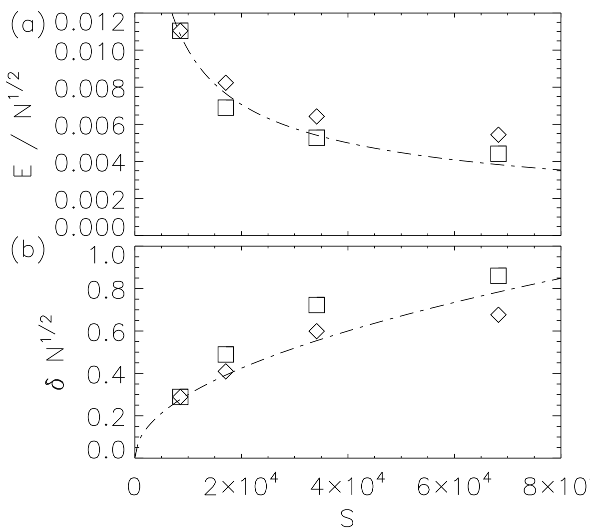

Fig. 2 Scaling of (**a**) reconnection rate E and (**b**) current layer thickness δ as a function of Lundquist number S in resistive-MHD simulations, from (Cassak et al. 2009), normalized to $N^{1/2}$ as in (2) and (3). The *dashed line* gives the prediction from the classical Sweet-Parker model. *Diamonds* and *squares* denote two techniques of measuring N

If the Sweet-Parker model describes each segment, the thickness δ of the segments scales as (Daughton et al. 2009b)

$$\delta \sim \frac{\delta_{\mathrm{SP}}}{\sqrt{N}}, \tag{2}$$

where δ_{SP} is the thickness predicted by classical Sweet-Parker theory in (1). Since $E \sim \delta/L$ from (1), the reconnection rate scales as (Cassak et al. 2009)

$$E \sim E_{\mathrm{SP}}\sqrt{N}, \tag{3}$$

where $E_{\mathrm{SP}} \sim S^{-1/2}$ is the classical Sweet-Parker rate. Since $N > 1$, secondary islands speed-up Sweet-Parker reconnection and make the current layers thinner, consistent with previous arguments (Shibata and Tanuma 2001).

These results were recently shown to be consistent with numerical simulations. Using PIC simulations, Daughton et al. (2009b) showed that the layer becomes thinner when secondary islands arise. Using resistive-MHD simulations, Cassak et al. (2009) did a scaling study for a range of Lundquist numbers up to $S = 6.8 \times 10^4$ and showed that (2) and (3) describe the data rather well, as shown in Fig. 2.

The next logical question is how the number of islands N scales with the Lundquist number S. The only way to study this numerically is to do extremely large simulations at high S. This has been pursued by Bhattacharjee et al. (2009), Huang and Bhattacharjee (2010), who did simulations up to $S \sim 10^6$. Their conclusion is that $N \sim S/S_{\mathrm{crit}}$, a linear scaling with S. Using (1) and (3), this predicts a reconnection rate of $E \sim 0.01$ independent of system parameters for large S. An alternate interpretation leading to the same result was recently presented (Uzdensky et al. 2010).

This result is very important as E is orders of magnitude faster than the classical Sweet-Parker prediction for large S, including inferred coronal parameters. It is only an order of magnitude slower than rates typically observed during Hall reconnection. This has important observational and theoretical consequences if it scales up to larger S, as will be discussed further in Sect. 6. However, it is very important to realize that the numerical simulations performed thus far go up to $S \sim 10^6$, which is large for numerical simulations but is 6 to 8 orders of magnitude smaller than coronal values. Thus, extrapolating to coronal Lundquist numbers is not obviously possible at this stage, as it could be the case that global effects are more important at larger S.

The other important caveat to take into account is that the simple theory in (2) and (3) assumes that all secondary islands are essentially active at the same time. However, it was proposed some time ago (Shibata and Tanuma 2001) that secondary island formation occurs hierarchically, in the sense that a global layer breaks in two due to a secondary island, then those two go unstable, and so on. This scenario was observed in resistive-MHD simulations (Bhattacharjee et al. 2009). This will have more of an effect for larger S, and will need to be incorporated into future theories.

Another result of importance is the role of embedded effects, which occur when the length scale over which the magnetic field changes direction on a global scale greatly exceeds the thickness of the dissipation region. This is expected to be the case in the corona, where dissipation takes place on the scale of meters, which is microscopic compared to global scales. The result is that the magnetic field that enters into the determination of the Lundquist number S is the magnitude of the reconnecting field immediately upstream of the dissipation layer (Cassak and Drake 2009). This is potentially far smaller than global characteristic field strengths, so it is important to keep this in mind when applying theoretical results to solar observations. Unfortunately, it is well beyond the capability of present instruments to measure local fields at reconnection sites at this time.

5.3 Impact of Secondary Islands on Onset

The other key element of secondary islands in Sweet-Parker reconnection that has recently emerged is its impact on the onset of collisionless reconnection. As discussed in Sect. 4, collisionless reconnection onsets abruptly when the layer reaches kinetic scales. It was proposed (Shibata and Tanuma 2001) that the formation of secondary islands makes the layer thinner (as in (2)), which could be a mechanism to usher in smaller length scales. Alternately, a similar effect occurs in embedded reconnection as the inflow convects in stronger reconnecting magnetic fields, which from (1) causes the layer to become thinner (Cassak et al. 2006).

The effect of secondary islands on the transition to collisionless reconnection was addressed numerically only recently. As mentioned in the previous subsection, it was observed in PIC simulations (Daughton et al. 2009b) that secondary island formation led to thinner current layers. These thinner layers became sub-gyroscale, and an onset of collisionless reconnection was observed. Again, these PIC simulations self-consistently calculate the effect of collisions, so it is possible to see that the reconnection electric field exceeds the Dreicer field after onset and collisionless effects become the dominant dissipation mechanism (Daughton et al. 2009a; Roytershteyn et al. 2010).

In these studies, however, numerical constraints forced S to be small enough that collisionless reconnection began as soon as a secondary island formed. It is important to separate whether the onset occurs as a result of the secondary islands or as a result of reaching kinetic scales. A recent resistive-Hall-MHD simulation study (Shepherd and Cassak 2010) showed that secondary islands can exist on scales larger than kinetic scales. The reconnection rates observed ($E \sim 0.01$) are limited to values consistent with the Bhattacharjee et al. (2009), Huang and Bhattacharjee (2010) studies. The faster reconnection rate of $E \sim 0.1$ was only observed to occur once the layers fell below kinetic scales and Hall reconnection began. Further, since the Hall reconnection rate is an order of magnitude faster than the secondary island reconnection rate as determined so far in simulations, it was observed (Shepherd and

Cassak 2010) that secondary islands were ejected from the vicinity of Hall reconnection sites once they became active. This implies that the energy release during the fastest phases of reconnection occur at Hall reconnection sites as opposed to at collisional reconnection sites with secondary islands.

One concludes from the studies discussed herein that (1) the condition that the layer needs to become smaller than kinetic scales in order to start collisionless reconnection seems to persist in the presence of secondary islands, (2) the presence of the secondary islands hastens the onset of collisionless reconnection, and (3) the most efficient energy release occurs during Hall reconnection, which dominates the reconnection process (if kinetic scales are reached).

6 Discussion

While the presence of secondary islands has been known for some time, the profound affect on the Sweet-Parker reconnection process was not broadly appreciated until recently. Recent work has furnished a better understanding of how reconnection proceeds in various parameter regimes (see Daughton and Roytershteyn 2011, this issue). Here, we summarize what has been learned and what open questions remain.

6.1 What Has Been Learned?

First, we consider the impact of what has been learned on the storage of energy in the build-up phase before a flare. In particular, we discuss recent suggestions that the pre-flare corona is undergoing reconnection, but it is collisional and therefore slow enough to allow energy to accumulate (Cassak et al. 2005; Uzdensky 2007; Cassak et al. 2008). As has been discussed earlier, there are three reconnection regimes: Sweet-Parker, Sweet-Parker with secondary islands, and collisionless (Hall) reconnection. There is evidence that reconnection in both the secondary island and collisionless phase are quite rapid ($E \sim 0.1$ for Hall reconnection, $E \sim 0.01$ for secondary island reconnection). If both of these scalings persist to large systems with high Lundquist numbers, then both are too fast to allow energy storage before a flare. Then, in order for pre-flare reconnection to be slow within the confines of these three models, it would have to be in the Sweet-Parker phase. To be in the Sweet-Parker phase, two criteria must be satisfied: (1) the thickness of the layer must be greater than kinetic scales, and (2) the Lundquist number must be below $S_{\mathrm{crit}} \sim 10^4$.

The solar corona, whether in active regions or in the quiet corona, has an extremely large Lundquist number $S \sim 10^{12}$ based on characteristic coronal parameters, which seemingly precludes Sweet-Parker reconnection from happening. However, one must be careful because the global Lundquist number is not the relevant quantity; rather, it is the Lundquist number based on the upstream reconnecting magnetic field. Since the global Lundquist number is eight orders of magnitude larger than S_{crit} at which secondary islands occur, the reconnecting magnetic field would have to be eight orders of magnitude smaller than the characteristic coronal field for embedded effects to allow for Sweet-Parker reconnection to occur. While outside the realm of measurement, it seems rather unlikely that this is the case for an extended period of time. Another possibility is that three-dimensional effects or effects not yet seen in simulations with S only as high as 10^6 play an important role in throttling islands. Thus, it is impossible at this time to definitively say whether classical Sweet-Parker reconnection without secondary islands can ever occur in the corona, but one would conclude based on present knowledge that it is unlikely.

If it is true that classical Sweet-Parker reconnection is not accessible, then the energy storage models are not tenable within the guise of the three canonical forms of reconnection. The conclusion would be that reconnection could not occur before a flare. In other words, the onset of any reconnection will rapidly lead to a fast energy release, so energy cannot accumulate while reconnection is occurring. In addition to the microphysical models, this result is important for several ideal-MHD models of coronal eruptions. In particular, some models (Forbes and Isenberg 1991; Antiochos et al. 1999) assume that MHD is valid for a long time and that thin current sheets form without significant energy release. Based on what has been learned from simulations of the microphysics, such current sheets would immediately reconnect rapidly. Thus, understanding what prevents reconnection from starting remains an open and critically important question. One mechanism of suppression of reconnection that has been discussed is the effect of line-tying (see *e.g.,* Delzanno and Finn 2008; Huang and Zweibel 2009). From the point of view of fundamental reconnection research, a better understanding of embedded effects during the reconnection between two flux tubes is necessary to address some of the open questions.

Second, consider the energy release phase of a flare. Many observations of solar eruptions reveal a distinct rise phase where the X-ray flux increases from noise followed by an eruptive phase. It was proposed (Shepherd and Cassak 2010) that the rise phase corresponds to reconnection with secondary islands, and the eruptive phase occurs when there is a catastrophic transition to Hall reconnection. The factor of ten difference in reconnection rates in the two phases is consistent with observed differences between the rise phase and eruptive phase seen in observations of flux emergence reconnection (Longcope et al. 2005) and implosion beneath a CME (Liu and Wang 2010). Much more work needs to be done to see if this simple model can explain the observations and whether the assumptions of the model make it applicable to the corona, such as whether the rise phase reconnection electric field is sub-Dreicer and whether the thickness of the layer during the secondary island phase is larger than kinetic scales.

It should be noted that the conclusions on energy storage and release may be very different in the chromosphere (Litvinenko 1999; Chae et al. 2003; Litvinenko and Chae 2009). Owing to the higher plasma density and smaller length scales, both the Lundquist number ($S \sim 10^8$) and the kinetic scales are smaller than in the corona. There are large uncertainties about the characteristic scales, so it is an open question which parameter regime the chromosphere falls in. It is entirely possible that reconnection obeys the Sweet-Parker model or, more likely, the modified theory allowing for secondary islands. However, the chromosphere is only partially ionized, so a full description of reconnection in these settings needs to include the interaction of the plasma with neutrals.

6.2 Open Questions

While much progress has been made on understanding collisional magnetic reconnection with secondary islands and its impact on the onset of Hall reconnection for coronal conditions, many questions remain unanswered. Here, we summarize two:

– *Effects of Turbulence*—Secondary islands in the form discussed in the present review occur self-consistently due to a secondary tearing instability of previously laminar fields. However, the presence of turbulence can effectively introduce secondary islands even without the secondary instability (Matthaeus and Lamkin 1985; Lazarian and Vishniac 1999). Thus, in order for the present results to be applicable in the corona, secondary tearing has to be more important than any ambient turbulence. Since the corona is a turbulent medium, it is not clear that this is the case. Some studies have investigated the relative effect of secondary islands and ambient turbulence (Smith et al. 2004; Loureiro et al. 2009; Skender and Lapenta 2010; Huang and Bhattacharjee 2010), but much further work is necessary to ascertain which effect is more important and under what conditions.

– *Three-dimensional effects*—The present discussion was completely based on evidence from two-dimensional theory and simulations. The landscape of the reconnection process is much different in three dimensions (Intrator et al. 2009). In particular, in the case of reconnection with a guide field (which is the expected naturally occurring situation in the corona), reconnection can occur at the symmetry line, but also at off symmetry axis locations referred to in the fusion community as rational surfaces. In such a system, secondary islands, which become flux ropes in three dimensions, can interact with each other in ways impossible in two-dimensions, including going around each other or becoming braided. Islands on multiple rational surfaces make the magnetic field in such regions stochastic (Borgogno et al. 2005). In addition, there are instabilities in three dimensions that are not present in two dimensions (Huba et al. 1977; Drake et al. 1994; Zhu and Winglee 1996; Drake et al. 1997; Büchner and Kuska 1999; Daughton 1999; Horiuchi and Sato 1999; Rogers et al. 2000; Lapenta and Knoll 2003; Karimabadi et al. 2003; Drake et al. 2003; Ricci et al. 2005; Yin et al. 2008) which potentially greatly change the behavior of reconnection. It will take much further work on this subject to understand how reconnection progresses in three dimensions.

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