MAGNETIC RECONNECTION: RECURSIVE CURRENT SHEET COLLAPSE TRIGGERED BY “IDEAL” TEARING

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ABSTRACT

We study, by means of MHD simulations, the onset and evolution of fast reconnection via the “ideal” tearing mode within a collapsing current sheet at high Lundquist numbers ($S \gg 10^4$). We first confirm that as the collapse proceeds, fast reconnection is triggered well before a Sweet–Parker-type configuration can form: during the linear stage, plasmoids rapidly grow in a few Alfvén times when the predicted “ideal” tearing threshold $S^{-1/3}$ is approached from above; after the linear phase of the initial instability, X-points collapse and reform nonlinearly. We show that these give rise to a hierarchy of tearing events repeating faster and faster on current sheets at ever smaller scales, corresponding to the triggering of “ideal” tearing at the renormalized Lundquist number. In resistive MHD, this process should end with the formation of sub-critical ($S \leq 10^4$) Sweet–Parker sheets at microscopic scales. We present a simple model describing the nonlinear recursive evolution that explains the timescale of the disruption of the initial sheet.

Key words: magnetic reconnection – magnetohydrodynamics (MHD) – Sun: corona

1. INTRODUCTION

Magnetic reconnection is thought to provide the pathway for energy release in solar flares and other phenomena where energy is accumulated in the magnetic fields and currents of ionized high-temperature plasmas. What has remained difficult to understand has been the triggering and speed of the process itself, which, given the extremely large Lundquist ($S$) and Reynolds ($R$) numbers, implies that currents must collapse to extremely small scales before anything close to realistic timescales are approached.

Indeed, an active region in the solar corona has a typical spatial scale $L \approx 10^9$ cm, a magnetic field $B \approx 50$ G, density $\rho \approx 10^9$ cm$^{-3}$ and a temperature $T \approx 10^6$ K, hence a macroscopic Lundquist number $S \approx 10^{13}$. A flare emits a total energy of about $10^{32}$ erg on a typical timescale of a few minutes, which cannot be explained by simple magnetic field diffusion or by a reconnecting instability occurring on the macroscopic scale itself. It has been suggested then that the reconnection trigger might only be provided by kinetic effects, beyond the (resistive) magnetohydrodynamic description of the plasma, or, alternatively, by the fast plasmoid instability of Sweet–Parker current sheets (Cassak & Drake 2013 and references therein).

It has been predicted recently, using linear stability analyses, that the tearing mode instability should grow on ideal timescales (“ideal” tearing, or IT) once the inverse aspect ratio $a/L$ of a current sheet reaches a scale of $a/L \sim S^{-1/3}$, preventing the formation of the paradigmatic Sweet–Parker current sheet (SP) for which $a/L \sim S^{-1/2}$, about 150 times smaller than $S^{-1/3}$ for $S \approx 10^{13}$ and even smaller for greater $S$. Even more interestingly, the (linear) asymptotic tearing instability in thin current sheets at arbitrary aspect ratios predicted that the maximum growth rate normalized to the macroscopic Alfvén time increases nonlinearly with the aspect ratio, $\gamma a^2 S \propto (L/a)^{3/2}$ (see also Equation (1) below): this suggests that collapse to the critical threshold thickness may provide the trigger for fast reconnection, where “fast” here is meant for dynamics occurring on the ideal timescale (Pucci & Velli 2014; D. Del Sarto 2015, private communication; Landi et al. 2015; Tenerani et al. 2015).

At the intermediate Lundquist numbers usual to simulations, say, $S \approx 10^3$–$10^5$, the presence of plasma flows into (inflows) and out of (outflows) the current sheet affects the scaling of the critical threshold for IT by inducing the formation of more elongated current sheets, that is, of layers having inverse aspect ratios $a/L \sim S^{-\alpha}$, with $\alpha_c > 1/3$ (M. Velli 2015, private communication). Therefore, with the Lundquist number not sufficiently large, the distinction between the IT scaling and the SP scaling might seem academic. Nevertheless, the IT current sheet instability, as is typical for the tearing mode, produces a quasi-singular inner layer $\delta$, whose inverse aspect ratio follows the SP scaling of $\delta_i/L \sim S^{-1/2}$. Such an inner layer is, however, non-stationary, with inflows and outflows increasing exponentially in time—their ratio incidentally does not follow the SP scaling, but rather $u_{in}/u_{out} \sim S^{-1/3}$—together with the amplitude of the reconnecting field: seen from this inner layer, one may interpret the dynamics in terms of the “embedded reconnection” scenario (Cassak & Drake 2009), leading to greater energy storage and more efficient dissipation at the extremely large $S$ values of the solar corona.

Here, we study the onset and evolution of the tearing instability within a single collapsing current sheet by means of resistive MHD simulations at $S = 10^6$. In the following paragraphs, we first show that the transition to a fast tearing mode instability takes place during the collapse when the predicted threshold inverse aspect ratio $a/L \sim S^{-1/3}$ is reached. Second, we show that the secondary current sheets formed nonlinearly give rise to recursive tearing instabilities at increasingly smaller scales and faster super-Alfvénic timescales, which may also be well described by the flow-modified “ideal” tearing criterion and instability. The nonlinear evolution leads to a complete disruption in a timescale estimated at 0.05% of the macroscopic Alfvén time once initial tearing is triggered, which we model here in a corrected version of the
fractal reconnection scenario first proposed by Shibata & Tanuma (2001). Though the two-dimensional instability of current sheets and its subsequent nonlinear evolution within resistive MHD have been studied before, it is shown here for the first time that the IT instability scenario, appropriately modified by the effect of the reconnection velocity flows, provides a quantitative description of the various stages in the evolution of a reconnecting current sheet, including hierarchical secondary island formation and disruption. Our results also provide a coherent framework through which previous simulations (Loureiro et al. 2005; Lapenta 2008; Bhattacharjee et al. 2009; Daughton et al. 2009) may be more completely understood.

2. TRIGGER OF FAST RECONNECTION: SCENARIO

Let us consider a (local) current sheet of inverse aspect ratio \( a/L \), as the one given in Equation (2), with Alfvén speed \( v_A \) and magnetic diffusivity \( \eta \), in a high \( S \approx L \eta / \eta \gg 1 \), constant density plasma. We assume, as in Uzdensky & Loureiro (2014), that the current sheet is collapsing via some external process on a timescale \( \tau_c \approx a (da/dt) \), which we take of the order of the ideal timescale, \( \tau_\eta \approx \tau_A \approx L/c_\eta \). We extend the linear analysis of the incompressible tearing instability to this time-dependent case by defining an instantaneous growth rate \( \gamma(k, t) \), \( k \) being the wave number along the sheet, where explicit time dependence is introduced by the evolving aspect ratio \( L/a \) itself.

In the static case, two regimes describe the unstable spectrum of a current sheet: the so-called small \( \Delta' \) regime, for wavelengths close to the instability threshold (or const.-\( \psi \), for \( ka \ll 1 \) for our equilibrium; Furth et al. 1963) and the large \( \Delta' \) regime (or non const.-\( \psi \), for \( ka \ll 1 \); Coppi et al. 1976). These two regimes have wave vectors \( k \) that lie to the right and to the left of the fastest growing mode \( k_{mn} \), \( k_m < k < 1/a \) and \( 0 < k < k_m \), respectively (Bhattacharjee et al. 2009; Loureiro et al. 2013; D. Del Sarto 2015, private communication). The wave vector and the growth rate of the fastest growing mode are given by (Pucci & Velli 2014)

\[
k_m a \sim S^{-1/4} / (L/a)^{1/4}, \quad \gamma_m \tau_A \sim S^{-1/2} / (L/a)^{3/2}.
\]

Similarly to Equation (1), the growth rates in the small and large \( \Delta' \) regimes also increase with the aspect ratio; however, their values tend to zero in the asymptotic limit \( S \rightarrow \infty \), and the full \( \gamma(k) \) dispersion relation becomes rapidly peaked around \( k_m \) (cf. Figure 1 of Pucci & Velli 2014). Therefore, both the small and large \( \Delta' \) regimes can be neglected in the framework of fast reconnection since the development of the instability is controlled by the evolution of the fastest growing mode described by Equation (1). Take the current sheet given by Equation (2), for example: the unstable spectrum lies within the range \( 2\pi/L < k < 1/a \). From this condition and the first of Equation (1), it follows that the regime described by Equation (1) always exists since \( k_m L \gtrsim 2\pi \), i.e., once \( a/L \gtrsim 0.2 S^{-1/5} \). The second of Equation (1) shows that the growth rate becomes ideal when approaching the critical width \( a/L \sim S^{-1/3} \) (much smaller than \( S^{-1/5} \)) from above: the transition to “ideal” tearing during collapse therefore occurs on the fastest growing mode. In conclusion, a collapse with \( \tau \gtrsim \tau_A \) naturally drives a sudden switch from a quasi-stable state—growth rate depending on a negative power of \( S \)—to an ideally unstable one. Otherwise, the disruption of the current sheet at widths thicker than critical (as considered in Uzdenskly & Loureiro 2014) would require an infinite time for \( S \rightarrow \infty \).

We now consider the example of an exponential collapse. Exponentially thinning current sheets are observed in simulations of solar and stellar coronal heating (Rappazzo & Parker 2013) in 3D MHD turbulence (Grauer & Marliani 2000; Brachet et al. 2013) of interest to a broad range of astrophysical phenomena, as well as in the nonlinear stage of the classic tearing instability itself (Ali et al. 2014).

3. NUMERICAL SETUP

We consider a plasma with homogeneous density \( \rho_0 \) and pressure \( p_0 \). The background magnetic field \( B_0 \) describes an exponentially shrinking current sheet and is given by

\[
B_0 = B_0 \tanh [(y/a(t))/\delta] + B_0 \sech [(y/a(t))/\delta],
\]

where the half-width \( a(t) \) is prescribed and parameterized in time by

\[
a(t) = a_0 \exp^{-t/\tau_L} + a_\infty (1 - \exp^{-t/\tau_F}).
\]

We employ a 2.5D compressible MHD code periodic in \( x \) and with non-reflecting boundary conditions along the inhomogeneous \( y \) direction (Landi et al. 2005). We assume an adiabatic closure, with index \( \Gamma = 5/3 \), a scalar resistivity, and a Newtonian viscous stress tensor with Prandtl number \( P = 1 \). With this choice, viscosity does not affect the scalings of IT significantly (Tenerani et al. 2015). The collapse of the background magnetic field is obtained by adding a source term \( \mathcal{F} \) of the form

\[
\mathcal{F} = \left( \frac{1}{a} \frac{\partial a}{\partial t} \right) y \frac{\partial B_0}{\partial y}
\]

in Faraday’s equation. Magnetic and velocity fields are normalized to \( B_0 \) and to \( v_A = B_0 / \sqrt{4\pi \rho_0} \), respectively, lengths to the macroscopic length \( L \), time to \( \tau_A \), and density and pressure to \( \rho_0 \) and \( B_0^2 / 4\pi \), respectively. It is also useful to introduce the normalized flux function \( \psi \), such that \( B = \nabla \times \psi \). We set \( p_0 = 0.8 \) and \( S = 10^6 \). Instability is seeded with a random noise of small amplitude. Because of the wide range of scales to be resolved in the \( y \) direction, we employ an inhomogeneous grid with increasing resolution at the neutral line. The simulation box has normalized dimensions \( L_x \times L_y = 2\pi \times 0.96 \) (runs 1–3) and \( L_x \times L_y = 2\pi \times 0.46 \) (run 4), with 2048 \( \times \) 1024 mesh points. Resolution at the neutral line is \( \Delta y = 0.0001 \), which allows one to resolve the diffusion region of the IT that scales as \( k_L/L \sim S^{-1/2} \) (Pucci & Velli 2014). Since for \( a \approx L \) the instability is extremely slow,

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**Table 1**

| Run # | \( S \) | \( \tau_c/\tau_A \) | \( a_{\infty}/L \) | \( \tau_\eta/\tau_A \) | \( \tau_{a}/\tau_A \) | \( a(\tau_a)/L \) |
|-------|-------|-----------------|-----------------|-----------------|-----------------|-----------------|
| run 1 | 1     | \( S^{-1/3} \) | 3               | 16              | 0.01            |
| run 2 | 10^6  | \( S^{-1/3} \) | 11.5            | 22              | 0.0105          |
| run 3 | 1     | \( S^{-1/3} \) | 16              | 36              | 0.013           |
| run 4 | 1     | \( S^{-1/3} \) | 4               | 4.5             | 0.0024          |
we start from $a_0 = 0.1 L$. This is a good compromise for limiting computational time while retaining a sufficient dynamical range for the collapse.

4. RESULTS

In Table 1, we list the background and main dynamical parameters for each run, and in Figure 1, we show for reference the temporal evolution of some unstable Fourier modes of theux function $\psi_k(y)$ in $y = 0$, from run 2. The growth rate and wavenumber of the fastest growing mode of the primary tearing at the critical IT threshold will be labeled $\gamma_i$ and $k_i$, respectively. In particular, in our simulations we find $k_i L = 10$ and $\gamma_i \tau_A = 0.46$, in agreement with theory (Tenerani et al. 2015).

4.1. Linear Stage

Tearing onset takes place at time $\tau_\text{nl}$ such that $\gamma_m(\tau_\text{nl}) \tau_\text{nl} = 1$, where $\gamma_m = 0.46 S^{-1/2} (a/L)^{3/2}$. After onset, the most unstable modes grow according to the WKB solution $\psi_m(t) \sim \exp \left[ \int \gamma_m(t') dt' \right]$, represented by the dashed line in Figure 1. In runs 1–3, we consider different cases of collapse in which $a_\text{\lower2pt\smallinfty}/L = S^{-1/3}$. In all these runs, the linear stage is ultimately dominated by modes close to the ideally unstable one, $k_i L = 10$. Run 4 forces a collapse of the initial sheet down to $a_\text{\lower2pt\smallinfty}/L = S^{-1/2}$, but proves that inverse aspect ratios $a/L \ll S^{-1/3}$ cannot be formed: the tearing mode indeed ignites right after $t = 2.4 \tau_A$, at which $a/L = S^{-1/3}$, and a large number of islands rapidly pop up. Even at the end of the linear stage (see Section 4.2), the collapsing sheet thickness is twice the corresponding SP thickness at the given $S$ and $P$ (Tenerani et al. 2015).

4.2. Early Nonlinear Stage

Nonlinearity becomes important at $\tau_\text{nl}$, when the width $w$ of magnetic islands is of the order of the width of the inner diffusion layer $\delta$. The fastest growing mode has $\delta_m$, which scales with the aspect ratio as $\delta_m/a \sim (a w / \eta)^{1/4}$ (Loureiro et al. 2013), hence $\delta_m/a \sim S^{-1/4} (L/a)^{1/4}$. By using the definition $w/a \simeq 2 \sqrt{\psi_m/a}$ (Biskamp 2000), we estimate the amplitude $\psi_m(\tau_\text{nl}) \simeq 0.25 S^{-1/2} \left[a(\tau_\text{nl})/L\right]^{1/2} \simeq 2.5 \times 10^{-5}$, in agreement with simulations. Nonlinear effects lead to the competition of two processes (Malara et al. 1991): on the one hand, islands start to merge through an inverse cascade process from the fastest growing modes to smaller wavenumbers, see, e.g., the inverse cascade from $k_i L = 10$ toward $k L = 5$ and below starting at $\tau_\text{nl} = 22 \tau_A$ (Figure 1); on the other hand, X-point collapse, and subsequent secondary current sheet formation, takes place during the further nonlinear growth of the islands, as is typical for strongly unstable modes far from the small $\Delta'$ regime (Jemella et al. 2003), in lieu of the slow, algebraic Rutherford growth (Rutherford 1973).

4.3. Nonlinear Stage: Recursive X-point Collapse

We analyze the fully nonlinear stage using run 1 as a reference, shown in Figure 2. X-point collapse leads to the formation of secondary plasmoids (magnetic islands; Loureiro et al. 2005; Lapenta 2008), and next, a recursive process of super-Alfvénic secondary layer formation and disruption takes place (Shibata & Tanuma 2001; Bhattacharjee et al. 2009; Daughton et al. 2009). Figure 2 shows a temporal sequence of recursive plasmoid formation, taking place at the center, near the flow stagnation point of the original instability. Previously, such recursive reconnection has been modeled as a succession of unstable SP layers (Loureiro et al. 2005, 2012; Cassak & Drake 2009; Daughton et al. 2009; Huang & Bhattacharjee 2010; Uzdensky et al. 2010). Our simulations show instead that it is driven by the onset of the IT mode, triggered by the dynamical lengthening of sheets to the local critical threshold, in a way similar to that discussed in Section 4.1. As the
reversive X-point collapse occurs in steps, we label the half-length of the \( n \)th current sheet \( L_n \), the inverse aspect ratio \( a_n/L_n \), the normalized (half-width) of the inner diffusion layer \( \delta_n/L_n \), the local Lundquist number \( S_n = (L_n/a_n) \tau_\alpha \), and the Alfvén time \( \tau_\alpha = (L_n/a_n) \tau_\alpha \). Referring to Figure 2, the reversal X-point collapse starts with the formation of \( L_1 \) from the primary diffusion region of thickness \( \delta_1 \), hence \( a_1 \approx \delta_1 \). Current sheet \( L_1 \) becomes unstable, less than 3 \( \tau_\alpha \) after the end of the linear stage (second panel), consistent with what was found in Landi et al. (2015). Here, we show that IT triggering within \( L_1 \) induces the growth of two plasmoids and of a second current \( L_2 \), with \( a_2 \approx \delta_1 \) (third panel), which is itself destroyed by multiple plasmoids in about 0.5 \( \tau_\alpha \) (fourth panel). In Figure 3, we plot \( B_y(y,0) \) at \( t = 0 \), \( t = 19 \tau_\alpha \), \( t = 19.8 \tau_\alpha \), the inset is a blow-up of \( B_y(y,0) \) at \( t = 19 \tau_\alpha \) and \( t = 19.8 \tau_\alpha \).

![Figure 3](image-url) Profile of \( B_y(y,0) \) in \( x_0/L = 3.2 \) at \( t = 18 \tau_\alpha \) (primary tearing, \( n = 0 \)), \( t = 19 \tau_\alpha \) (\( n = 1 \)), and \( t = 19.8 \tau_\alpha \) (\( n = 2 \)). The inset is a blow-up of \( B_y(y,0) \) at \( t = 19 \tau_\alpha \) and \( t = 19.8 \tau_\alpha \).

![Figure 4](image-url) Length \( L/L \) (green), width \( a/L \) (magenta), and inverse aspect ratio \( a/L \) (blue) of the first secondary current sheet vs. time (run 1). The red dotted line represents the critical threshold \( S^{-1/3} \), and the light blue dotted line the viscous \( S_{visc} \). Black dotted and dashed-dotted lines represent the critical threshold with outflow effects, given by Equation (5).

a threshold factor \( f \approx 0.5-0.1 \). The value \( \mu = 10 \) (Biskamp 2000) yields the observed \( \alpha_c = 1/2 \) for \( S_n = 10^4 \), while, as expected, \( \lim \alpha_c = 1/3 \). Note that, as shown below, \( S_n \), is a decreasing function of \( n \). The black dotted and dashed-dotted lines in Figure 4 correspond to \( S_n^{1-\alpha_c} \) at two different \( \mu \), while the light blue dotted line corresponds to \( a_{visc}/L \). Though our Lundquist numbers are not extremely large, \( L_1 \) still disrupts before reaching the SP width. The trend of the data plotted in Figure 4 goes in the direction of our scenario, that is, that the fast tearing instability is triggered within collapsing current sheets once the critical threshold is reached. Taking into account the increased Alfvén speed due to pile-up just outside the inner diffusion region as considered by Cassak & Drake (2009) does not change the trend of our results. Interestingly, the discussion of flux pile-up and embedded reconnection led (Cassak & Drake 2009) to independently finding a similar scaling \( \sim S^{-1/3} \) for the inverse aspect ratio of unstable current sheets starting from an initial secondary Sweet–Parker sheet. Their analysis starts from a different line of thought in search of the proper criterion to destabilize an embedded SP sheet and proceeds via a linearization of the fields in the neighborhood of the neutral line. It then finds a scaling with \( S \) for the aspect ratio of the embedded layer scaling in a way similar to IT. On the other hand, the IT scaling derives from requiring an \( S \)-independent growth rate from the complete eigenmode analysis of tearing instability theory and finds that the entire sheet inverse aspect ratio scales as \( S^{-1/3} \). The results therefore have different origins, and for the IT case, the presence of flows modifies the scaling exponent \( \alpha_c \) according to Equation (5); incidentally, the latter expression is in excellent agreement with the aspect ratios at which plasmoids are observed to be ejected in the Cassak & Drake (2009) paper.

5. DISCUSSION

We have shown that a collapsing current sheet disrupts due to the triggering of fast reconnection as predicted by the IT model. This process is shown to proceed recursively, giving

In Equation (5), \( \mu = \Gamma / (\Omega \beta) \) is the ratio of the plasmoid evacuation rate to the maximum growth rate, corrected by

\[
\alpha_c = \frac{2 \log \mu + \log S_n}{3 \log S_n}.
\]
rise to a hierarchy of fast tearing instabilities at ever smaller scales. Let us model the recursive collapse of X-points, in the spirit of the original fractal reconnection model, with the following differences: first, we do not consider SP as the initial condition; second, we use the width $a_n$, derived from the IT scalings, to find $L_n$. Neglecting for simplicity viscous and outflow effects, imposing that $a_n \ll L_n$, we find $n \to \infty$, then $\tau_{A,n} \to \tau_A S^{-1}$, and, as expected, $S_n \to 1$ and $a_n/L_n \to 1$. However, after a number $n_*$ of steps, $S_{n_*} \approx 10^{-4}$, at which we expect to reach a stable SP. For $S = 10^{13}$ (solar corona), $n_* \approx 4$, and the length $L_n$ suddenly drops to microscopic scales, $L_4/L \approx 6 \times 10^{-4}$, $L_2/L \approx 2 \times 10^{-6}$, $L_3/L \approx 3 \times 10^{-8}$, and $L_4/L \approx 10^{-9}$. The recursive X-point collapse after the first trigger thus lasts a time interval of $\tau_\text{tot} \approx \sum_{n=1}^{n_*} \tau_{A,n} \approx 5 \times 10^{-4} \tau_A$, one order of magnitude less than the upper limit given by Shibata & Tanuma (2001).

Higher-resolution simulations are nevertheless necessary to assess how the nonlinear evolution saturates. Our scenario can be modified to include kinetic scales that might be reached dynamically in typical astrophysical systems (Cassak et al. 2005; Daughton et al. 2009). In this case, we expect a change of the power laws (6) as different scalings describe the physics at the critical threshold in the kinetic regime (D. Del Sarto 2015, private communication).

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