IUTAM Symposium on Understanding Common Aspects of Extreme Events in Fluids

On the relaxation of braided magnetic fields

D. I. Pontin*, A. L. Wilmot-Smith, G. Hornig

Division of Mathematics, University of Dundee, Dundee, DD1 4HN, U.K.

Abstract

We discuss the properties of braided magnetic fields, in particular with regard to their relaxation towards equilibrium both in the absence and presence of finite plasma resistivity. We consider a model magnetic field in which a subset of field lines are braided in the mathematical sense (in the form of a pigtail braid), and stretch between two perfectly conducting plates. An ideal relaxation yields an approximate force-free equilibrium without thin current sheets. However, in the presence of a finite resistivity the field spontaneously develops thin (but finite) current layers. Subsequently, the system undergoes a fast relaxation on an Alfvén timescale via a myriad of thin current layers. The current layer structure becomes increasingly complex, and is increasingly long-lived, for higher magnetic Reynolds numbers. The effect of the resistive relaxation is to unbraid the magnetic field, that is to significantly simplify the mapping of field lines between the end plates. This process involves magnetic flux being reconnected multiple times. The final state of the resistive relaxation approximates a non-linear force-free field, implying the presence of additional constraints beyond the magnetic helicity on the relaxation process. The ‘topological dissipation’ hypothesis, which proposes to explain the heating of the solar corona via braiding of the coronal field by photospheric motions, is one of the most controversial ideas of solar physics. We therefore consider in some detail the relevance of our results to this debate.

Keywords: Magnetohydrodynamics; magnetic topology; magnetic reconnection; solar corona

1. Introduction

In this paper we study the properties of braided magnetic fields. The original motivation for this study was to understand the so-called topological dissipation mechanism for heating the solar corona. This stems from the observation that the plasma of the solar corona is at a temperature of order $10^6 K$, while the solar surface temperature is only around $10^4 K$. The plasma-$\beta$ in the corona is small, meaning that magnetic forces dominate the dynamics. By contrast, $\beta$ is large in the solar interior, so that the flow dominates. Due to the turbulent convection in the outer layer of the solar interior, the foot points of coronal magnetic field lines are therefore continually shuffled around. Since the magnetic Reynolds number, $R_m$, in the corona is high, the photospheric motions tangle the coronal magnetic field lines around one another, and it was Parker [1] who first proposed that the associated excess energy stored in the coronal field could be liberated as a source of plasma heating.

*Corresponding author
Email address: dpontin@maths.dundee.ac.uk (D. I. Pontin)
In order to release all of the energy associated with the braiding or tangling of the coronal magnetic field, one requires to simplify the topology. This can only happen as a result of magnetic reconnection [for reviews see e.g. 2, 3, 4]. However, owing to the high value of $R_m$, reconnection may only take place if sufficiently intense current layers form. Now, since the Alfvén travel time along a typical coronal loop is much shorter than the timescales associated with the turbulent photospheric convection, it is usually thought that the coronal magnetic field should evolve through a sequence of equilibrium states (approximate force-free states satisfying $\mathbf{J} \times \mathbf{B} = 0$ since $\beta \ll 1$). In addition, since the plasma is to a good approximation ideal, the relaxation towards each of these equilibrium states should preserve the topology of the coronal field. Parker and others [1, 5, 6, 7, 8] have argued that relaxation to a smooth force-free state following an arbitrary perturbation of the field is impossible (except in artificial cases with certain symmetries). Instead they propose that tangential discontinuities, corresponding to singular current sheets, would develop in the field if the relaxation were truly ideal. It is thus proposed that in the presence of even a very small finite resistivity, current layers develop that are sufficiently intense to allow reconnection to occur. Once the topology of the field lines is allowed to change, the magnetic field may seek a lower energy state, releasing additional magnetic energy.

The topological dissipation mechanism continues to generate much debate, and in particular a number of authors have argued against the development of singular currents in the ideal limit [9, 10, 11, 12, 13, 14]. In the braided magnetic fields considered herein, we find no indication of the formation of singular current layers. However, we propose that braiding of the magnetic field is indeed a possible mechanism for heating the corona, albeit in a rather different way to that initially proposed. It is to be anticipated that braided magnetic fields exist in many systems other than the solar corona. As such, we describe here the relaxation of such braided fields in general terms. In Sections 2 and 3 we introduce the model and the numerical methods employed. In Sections 4–6 we describe the results and in Section 7 present some conclusions.

2. The model magnetic field

The model magnetic field we consider is contained within a volume in Cartesian geometry in which magnetic flux enters the volume through $z = -z_0$ ($z_0$ constant) and leaves through $z = z_0$ (there are no closed field lines or null points; $B_z > 0$ throughout). Regarding the application to magnetic loops in the solar corona we are neglecting the effects of the curvature of coronal magnetic field lines; the planes $z = \pm z_0$ can be considered as separate regions of the solar photosphere. We take here $z_0 = 24$. Owing to the restricted range of $R_m$ accessible in numerical simulations, we choose to begin our study with an initial magnetic field in which the field lines are already braided. For simplicity we confine our discussion here to one particular initial braided state, modelled on the pigtail braid (see Fig. 1(a)). In an idealised cartoon picture, such a braid can be built up from three initially straight strands by sequentially exchanging adjacent pairs of footprints at one boundary. This could be achieved using two opposite-sign rotational motions, denoted $\sigma_1$ and $\sigma_1^{-1}$ in Fig. 1(a), where each represents a rotation by $\pm \pi$ radians. Performed in the sequence $\sigma_1, \sigma_1^{-1}, \sigma_1, \sigma_1^{-1}, \sigma_1, \sigma_1^{-1}$, this procedure generates the pigtail braid.

Of course a magnetic field is not made of discrete strands but is space filling. We construct our initial magnetic field as follows. We take a uniform ‘background’ magnetic field $B_0 e_z$, and superimpose upon it six localised regions of closed magnetic flux corresponding to regions of localised electric current. Each of these perturbations is a flux ring lying in the $xy$-plane (with $B_z = 0$), and so when added to the vertical background field acts to create a local twist of magnetic field lines, though a maximum angle of around $\pi$ radians – see Fig. 1(b). Specifically, we take

$$
\mathbf{B} = B_0 \left[ e_z + \sqrt{2} \sum_{i=1}^{6} k_i \exp \left( -\frac{(x-x_i)^2 + y^2}{2} + \frac{(z-z_i)^2}{4} \right)(-y e_x + (x-x_i)e_y) \right]. \tag{1}
$$

setting $B_0 = 1$, and choose three regions of positive twist centred on $x = 1, y = 0, z = \{-20, -4, 12\}$ and three regions of negative twist centred at $x = -1, y = 0, z = \{-12, 4, 20\}$. This is obtained by setting $k_{1,2,3} = 1$; $x_{1,2,3} = 1$; $z_{1,2,3} = \{-20, -4, 12\}$, and $k_{4,5,6} = -1$, $x_{4,5,6} = -1$, $z_{4,5,6} = \{-12, 4, 20\}$. The magnetic field defined in this way has field lines that are linked in a range of different ways. Importantly, a subset of the
field lines in the domain are braided around one another with the pigtail braid structure of Fig. 1(a). The braiding induced in the field can be visualised by plotting the mapping of field lines from $z = -24$ to $z = 24$. In Fig. 2 we plot the squashing factor, $Q$ [15], associated with this mapping, in which a complex pattern of thin layers is present [see also 16].

It is worth noting a number of properties of this magnetic field that are relevant to the coronal heating problem. First, the aspect ratio of our model coronal loop is relatively realistic when compared to previous computational studies of flux braiding in the corona, which often considered very short (or wide) loops [17, 18, 19]. More specifically, previous studies often applied shear via boundary driving that resulted in a displacement (perpendicular to the loop) of field line footpoints between the end plates of similar size to the loop length. By contrast, the maximum footpoint displacement of field lines between the end plates in our configuration is an order of magnitude smaller than the loop length. The low level of twist of field lines in our model field is also consistent with observations, in which typically loops do not appear highly twisted on large scales. One further point to note is that the perturbation applied to the background field by adding the regions of twist is not a violent one. Indeed, the addition of the perturbations only increases the net energy of the field above that of the background homogeneous field by around 5% (within $x, y \in [-3, 3]$, $z \in [-24, 24]$ where the dynamics described below are approximately confined).

3. Numerical methods

The magnetic field as described in the previous section is taken as an initial condition for numerical simulations. These can be described as “relaxation” simulations, since the initial field is not an equilibrium, and we do not perform any driving of the system at any time. The simulations are run in a computational domain defined by $x, y \in [-6, 6]$, $z \in [-24, 24]$. On all boundaries the magnetic field is line-tied (field line footpoints are fixed), and the plasma velocity is set to zero. The magnetic field is essentially parallel to the $x$-
and $y$-boundaries, which do not significantly influence the simulations (the dynamics are confined roughly within $x, y \in [-3, 3]$).

There are two separate stages to the relaxation procedure. In the first stage, we perform an ideal relaxation using a Lagrangian magneto-frictional approach that precisely preserves the topology of the magnetic field [for details see 20, 21]. The final state of the ideal relaxation is a numerical approximation to a force-free field (satisfying $\mathbf{J} \times \mathbf{B} = 0$) – the accuracy of this approximation is discussed below.

In the second stage of the relaxation process we introduce a small but finite plasma resistivity, and solve the full system of resistive MHD equations to follow the associated resistive relaxation. The simulations are run with the finite-difference-based 

\textit{Copenhagen StaggerCode}, see [22]. We employ a prescribed spatially uniform resistivity and a combined second- and fourth-order hyper-viscosity model that switches on around shock structures. The numerical resolution used is $512^3$. At the beginning of this resistive relaxation phase we reset the plasma density and pressure to spatially uniform values ($\rho = 1, p = 0.067$) to remove any pressure gradient forces. The result is a plasma-$\beta$ that varies throughout the domain between 0.10 and 0.14 at $t = 0$. Further details of the simulations may be found in [23, 24].

4. Results: Ideal relaxation process

The ideal Lagrangian relaxation scheme described in the above section is run with our braided magnetic field (see Fig. 1(b)) as initial condition, using $161^3$ grid points. In the final state of this ideal relaxation, the twist associated with the perturbations inserted into the field has equilibrated along the $z$-direction – see Fig. 1(c). This corresponds to a smoothing of the current along the $z$-direction, as shown in Fig. 3(a). The current structure in the relaxed state has only large-scales, i.e. no thin current layers are found. In addition, the excess magnetic energy is reduced from 5% to 1.7% above that of the background homogeneous field. For further details of this ideal relaxation, see [21].

The final state of the ideal relaxation is of course only a numerical approximation to a force-free field. The relaxation is limited by roundoff errors that accumulate as the mesh distortion increases [25]. We estimate the proximity to a force-free field by evaluating $\mathbf{J} \times \mathbf{B}$ on the final numerical mesh using a \textit{mimetic} approach – the method is described in [25]. While the peak value of $\mathbf{J} \times \mathbf{B}/B^2$ at the start of the ideal relaxation is 1.0, the peak value at the end is $2 \times 10^{-2}$.

In the following Section, we will describe the results of resistive MHD simulations on a high-resolution Eulerian grid in which the final state of the magnetofrictional (MF) relaxation is used as an initial condition. We will see that the field is unstable, and collapses to form a pair of thin current layers (of finite width). It
is worth noting, however, that even if the resistivity is set to zero in this Eulerian MHD code, we still find a collapse to form thin current layers. This leads us to question why no thin current layers are present in the approximate equilibrium obtained during the ideal MF relaxation. A number of possible explanations for this exist:

1. Thin (perhaps singular) current layers do exist in an exactly force-free state with the same topology as our braided field, but the MF relaxation cannot proceed to sufficiently low $\mathbf{J} \times \mathbf{B}$ to obtain the thin current layers, owing to the limitations arising from mesh distortion [25]. We note that the indications from the Eulerian MHD simulations with $\eta = 0$ are that any such thin current layers are of finite thickness.

2. The end state of the MF relaxation is close to an ideal unstable equilibrium, but either (a) the equilibrium is linearly stable, and only non-linearly unstable, or (b) the instability is linear but has not been seen in the simulations as it grows extremely slowly.

3. The instability is resistive, and even a small numerical resistivity in the MHD code (the code uses sixth-order spatial differencing) is sufficient to trigger it.

4. The end state of the MF relaxation is close to a stable equilibrium, but the perturbation associated with interpolating to the rectangular mesh required for the MHD code is sufficient to disturb this.

At present we do not know which of the above statements is true. This is a topic of current investigation.

5. Results: Resistive relaxation process

5.1. Relaxation for $\eta = 10^{-3}$

We now describe the results of switching on a small but finite resistivity, using the final state of the ideal MF relaxation as initial condition. We first focus on a simulation run using $\eta = 10^{-3}$. We find that the initially smooth structure of the current density is quickly destroyed through the formation of a pair of intense current layers after around 10 time units, see Fig. 3(b). (Note that the Alfvén travel time $\tau_A \approx 1$, and the perturbations to the field shown in Fig. 1(b) are 8 space units apart). This is associated with around a three-fold increase in the peak current modulus in the domain, as shown in Fig. 4. These current layers eventually diffuse away, but are replaced by an increasingly fragmented array of current layers that rather effectively fill the part of the domain containing braided magnetic flux, see Fig. 3(c, d). After $t \approx 70$, the
fragmentation of the system of current layers starts to reverse, with the structure of the current gradually simplifying. We note that the most intense current layers that form throughout the simulation are the two initial current layers (see Fig. 4); the peak current in the domain fluctuates rather rapidly, but shows a general tendency to decrease in time towards a final minimum value.

Since the resistive relaxation prior to \( t \approx 200 \) is facilitated by the formation of myriad thin current layers, it is a fast relaxation – the whole process occurs over a period of time equivalent to approximately 4 times the Alfvén transit time along the box in the \( z \)-direction. By \( t \approx 200 \), the system reaches a state with two large-scale current features that are extended along the \( z \)-direction (Fig. 3(e)). In the shortest direction through these structures they have dimension on the order of the system size (in \( xy \)). As such, after this time the continuing relaxation of the field through the dissipation of these currents will be slow, being governed by the timescale for resistive diffusion. While this timescale is only of order \( 10^3 \) non-dimensional time units for the present simulation parameters (only around 5 times longer than the period of fast relaxation), in a magnetic loop in the solar corona the timescale for resistive diffusion is enormous (order \( 10^{12} \) s), and so we consider the state reached after around \( t = 200 \) to be the physically relevant final state. To be conservative here we take for our final state \( t = 290 \). The nature of this final state is discussed in more detail in Section 6.

The presence of a finite resistivity allows the topology of the magnetic field to change, i.e. the mapping of field lines from \( z = -24 \) to \( z = 24 \) is not fixed. This change of topology is shown in Fig. 5, where we plot at different times four sets of field lines, two sets from fixed footpoints at \( z = -24 \) (green, blue) and two from fixed footpoints at \( z = 24 \) (red, orange). The field lines are selected in such a way that at \( t = 290 \) they become coincident. It is clear from the figure that there is an unbraiding of the magnetic field – the connectivity of field lines between \( z = \pm 24 \) is gradually simplified.

In order to quantify the changing topology of the magnetic field, we evaluate the magnetic reconnection rate throughout the simulation. Three-dimensional magnetic reconnection occurs in localised regions where the electric field has a component parallel to the magnetic field, \( E_\parallel \). The reconnection rate is defined for such a region [26, 27] by the maximal value of

\[
\Phi = \int E_\parallel dl
\]

over all magnetic field lines threading the region, where the integral is performed along magnetic field lines from one side of the non-ideal region (region within which \( E_\parallel \neq 0 \)) to the other. During the simulation, there are many 3D-localised non-ideal regions, corresponding to the many current layers discussed above (since \( E_\parallel = \eta J_\parallel \)). We therefore have to track many isolated non-ideal regions, and evaluate the reconnection rate for each individually. This is done by evaluating \( \Phi \) along all field lines and seeking localised maxima and minima – for further details of the procedure used see [24]. The global reconnection rate for the system is then obtained by summing the moduli of the local maxima and minima of \( \Phi \).

The results of performing the above analysis are shown by the dashed line in Fig. 6(a). We see that the reconnection rate, like the peak current, peaks sharply at early time before gradually decreasing towards the end of the simulation. Note, however, that the maximum global reconnection rate occurs much later than the maximum peak current. This is due to the fact that an increasing number of individual non-ideal regions (or reconnection sites) form as the fragmented current structures becomes increasingly complex (see the dashed

![Fig. 4. Spatial maximum of the current density modulus |J| as a function of time.](image-url)
Fig. 5. Selected magnetic field lines traced from fixed footpoint locations on $z = 24$ (red, orange) and $z = -24$ (green, blue), at (a) $t = 0$, (b) $t = 40$, (c) $t = 80$, and (d) $t = 200$.

Fig. 6. (a) Reconnection rate, (b) number of identified reconnection regions in the positive half-space $z > 0$, and (c) cumulative reconnected flux, for runs with $\eta = 10^{-2}$ (grey solid line), $\eta = 10^{-3}$ (black, dashed) and $\eta = 2 \times 10^{-4}$ (black, solid). After [24].
Fig. 7. Current density in the $z = 0$ plane at $t = 135$ for simulations with (a) $\eta = 10^{-3}$ and (b) $\eta = 2 \times 10^{-4}$.

line in Fig. 6(b)). The peak reconnection rate occurs at $t = 50$, when the largest number of reconnection sites is identified.

An interesting property of the system can be observed by evaluating the total magnetic flux that is reconnected during the relaxation process. This is obtained from a time-integral of the reconnection rate, and the cumulative reconnected flux during the simulation is shown by the dashed line in Fig. 6(c). The total flux reconnected by $t = 290$ is 41.2 units. This figure should be compared with the flux associated with the initial perturbation (or twist) inserted in the field, and the twist remaining in the final state (which must be non-zero owing to the currents present). The initial ‘poloidal’ flux is known from the analytical expression to be 30 units (flux through the $y = 0$ plane), and the high degree of symmetry in the final state allows the poloidal flux remaining to be estimated by evaluating the flux through the $y = 0$ plane, which is found to be 15.3 units. Thus, beginning with 30 units of flux, we reconnect 41.2 units, and end up with 15.3 units. The inescapable conclusion is that some of the magnetic flux is reconnected multiple times. Indeed, we see that on average each unit of flux must be reconnected approximately 2.8 times. This demonstrates the complexity of the topological change required for the magnetic field to unbraid itself during this resistive relaxation. Such multiple reconnection of flux has also been observed in simulations of the ‘fly-by’ of two flux sources, in which multiple separators are present [28]. We therefore think it likely that this is a generic feature of complex interactions in 3D magnetic fields.

5.2. Dependence on Reynolds number

Above we described the relaxation of our braided magnetic field for a resistivity of $\eta = 10^{-3}$. We have also repeated our resistive MHD simulations using the same initial condition but with $\eta = 10^{-2}$ and $\eta = 2 \times 10^{-4}$ (the lowest value that we can use with the present numerical resolution while still resolving the current layers that develop in the system). The differences between the simulations can be summarised as follows. For larger Reynolds number $R_m$:

1. There is a greater complexity in the current structure, i.e. more individual current layers and therefore more isolated reconnection sites, see Figs. 6(b), 7.
2. The relaxation process takes longer, i.e. it takes longer for the state with many current layers to eventually relax to the state with only large-scale current structures, see Figs. 6(a, b), 7.
3. The maximum current is larger, but the maximum reconnection rate for any single non-ideal region is smaller. However, the global reconnection rate depends at most weakly on $R_m$, see Fig. 6(a), due to the increasing number of reconnection sites.
4. The average number of reconnections per unit of flux required to unbraid the field increases, see Fig. 6(c).
6. Results: Final state of the resistive relaxation process

We lastly consider the properties of the final state in the resistive relaxation described above. As already discussed, the topology of the final state is rather simple compared with the initial state. In broad terms it consists of two twisted flux tubes embedded within the uniform background field in the $z$-direction. The two twisted flux tubes correspond to the two current concentrations evident in Fig. 3(e), and have opposite sense of twist (the current is dominated by the $J_z$ component, which has opposite sign in each current concentration). This structure can be visualised by plotting the mean value along field lines of the quantity $\alpha^* = J \cdot B / B \cdot B$.

Note that in a true force-free field satisfying $J \times B = 0$, to which the initial and final states are approximations, this quantity would be constant along field lines. In Fig. 8, $\alpha^*$ is plotted as a function of the intersection point of field lines with the $z = -24$ boundary for the central portion of the domain in $xy$. In the initial state the complex field line mapping is evident through the numerous thin layers present. By contrast we see that in the final state there are two large-scale regions of oppositely signed $\alpha^*$ corresponding to the flux tubes with positive and negative twist. Note also that the maximum and minimum values of $\alpha^*$ are reduced in the final state.

It is clear from Fig. 8 that the final state of the relaxation approximates a *non-linear* force-free field. This is inconsistent with the hypothesis of a Taylor relaxation, which was first put forward to explain turbulent relaxation in a reversed-field pinch [29], and was subsequently proposed as being applicable in the solar corona [30]. The theory predicts that the only constraint on a sufficiently turbulent relaxation is the total magnetic helicity ($H = \int A \cdot B \, dV$). However, the helicity in our system is exactly zero at all times due to the symmetry of the initial configuration, and a Taylor relaxation would therefore predict the uniform vertical magnetic field as the final state. We do not believe that the explanation for this failure of the Taylor prediction is an insufficient level of turbulence in the system, since for larger $R_m$ the relaxation is significantly more turbulent, and yet there is more twist remaining in the field in the final state (larger maximum and minimum values of $\alpha^*$). Instead, it has been discovered that an additional constraint on the relaxation is present, which prevents the relaxation to the homogeneous field. This is the topological degree of the field line mapping, given by the sum of the indices of fixed points of the mapping. This is discussed in more detail in the article by Yeates *et al.* in this volume, and in [31].

7. Discussion and conclusions

We have discussed the properties of a particular braided magnetic field, containing zero net helicity and modelled on the pigtail braid. It should be noted however that the qualitative properties discussed have been...
observed during the relaxation of a number of different braided magnetic fields [31, 32, 33]. One of the key points that remains to be understood is the existence of a smooth, equilibrium field with such a braided topology. While the numerical Lagrangian magnetofrictional relaxation code used herein yields a smooth approximate equilibrium, it can never reach an exact equilibrium. What is clear is that the (approximate) equilibrium obtained after the ideal relaxation is susceptible to some form of instability. This may be a resistive instability, as proposed in [21], although the fact that it sets in at the same time when the explicit resistivity is set to zero in our MHD code (leaving only numerical resistivity which is small due to the sixth-order spatial differencing) casts some doubt on this.

At the beginning of the resistive evolution, a collapse to two thin current layers is observed. This is followed by an increasingly complex system of current layers, that we believe would develop into a turbulent state were the magnetic Reynolds number sufficiently high. During this stage of the evolution, intense (but finite) current layers fill the volume rather effectively. Each of these localised current layers corresponds to a magnetic reconnection process, through which the topology of the field (i.e. the mapping of field lines from one end of the domain to the other) is changed. The net effect of these many isolated reconnection events is to unbraid the magnetic field, such that at the end of this fast reconnection-mediated resistive relaxation the field line mapping is greatly simplified. During this unbraiding of the field, it turns out that each unit of flux is reconnected multiple times (on average). The final state of our simulation contains two twisted flux tubes, and as such corresponds to an approximate non-linear force free field. This has led to the discovery of an additional constraint – the topological degree of the field line mapping – on the relaxation of magnetic fields of this geometry.

Returning to Parker’s topological dissipation hypothesis, our results suggest that the braiding of the coronal magnetic field may indeed be a viable mechanism to heat the coronal plasma. The details of how this is realised, however, appear to be different to those proposed by Parker. Specifically, in the mathematical idealisation of infinite magnetic Reynolds number, we find no singular current sheets (based on a numerical ideal relaxation that exactly preserves the topology of the magnetic field, but that achieves a force-free state only to a certain accuracy, as discussed above). However, we do observe a fast relaxation involving many thin, intense current layers for finite $R_m$, through which excess magnetic energy is liberated. In order to understand in detail the implications of magnetic braiding for the dynamics of the solar corona – or indeed other plasma environments – we need to understand how the relaxation depends on $R_m$, which in our resistive MHD simulations is many orders of magnitude smaller than the true value in the corona. We have performed simulations with a limited range of $R_m$, and find that for larger $R_m$ the system of current layers that develops becomes more complex and persists for longer, with the average number of reconnections per unit of flux increasing. The implications of our results for coronal dynamics and heating also require a more careful treatment of the plasma response to the heat that is deposited during the relaxation. This will require in future that energy losses are properly included in the energy equation solved in the simulations, and for example the effects of atmospheric stratification included. Moreover, thus far we have not treated the process that braids the field, instead assuming that the braiding has already been built up in the field. We need to understand how the braiding builds up self-consistently through boundary driving, and how this dissipated driving eventually leads to a balance between the energy injected by the driving and the energy dissipated through the reconnection events in the loop. This is a topic of current study.

References


