Self-Regulation of Solar Coronal Heating Process via the Collisionless Reconnection Condition

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I propose a new paradigm for solar coronal heating viewed as a self-regulating process keeping the plasma marginally collisionless. The mechanism is based on the coupling between two effects. First, coronal density controls the plasma collisionality and hence the transition between the slow collisional Sweet-Parker and the fast collisionless reconnection regimes. In turn, coronal energy release leads to chromospheric evaporation, increasing the density and thus inhibiting subsequent reconnection of the newly reconnected loops. As a result, statistically, the density fluctuates around some critical level, comparable to that observed in the corona. In the long run, coronal heating can be represented by repeating cycles of fast reconnection events (nanoflares), evaporation episodes, and long periods of slow magnetic stress buildup and radiative cooling of the coronal plasma.

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This Letter is devoted to the problem of solar coronal heating (see Ref. [1] for a recent review), viewed in the context of Parker's nanoflare model [2]. Since the main heating process in this model is magnetic reconnection, I will first summarize the recent progress in reconnection research achieved in the past 20 years. Even though a complete picture of reconnection is still not available, there is now consensus about some of its most important aspects. My main goal is to apply this new knowledge to the coronal heating problem.

First, I want to emphasize the importance of a realization by Petschek [3] that the main bottleneck in the classical Sweet-Parker [4,5] reconnection model is the need to have a reconnection layer that is both thin enough for the resistivity to be important and thick enough for the plasma to be able to flow out. Furthermore, Petschek [3] proposed that this difficulty can be mitigated if the reconnection region has a certain substructure—the Petschek configuration, with four shocks attached to a central diffusion region. This results in an additional geometric factor leading to faster reconnection. This idea is especially important for astrophysical systems, including the solar corona, irrespective of the actual microphysics inside the layer. Indeed, the system size L is usually much larger than any microscopic physical scale δ, e.g., the ion gyroradius ρ_i, the ion collisionless skin depth $d_i \equiv c/\omega_{pi}$, or the Sweet-Parker layer thickness $\delta_{SP} = \sqrt{L\eta/V_A}$. Therefore, a simple Sweet-Parker-like analysis would give a reconnection rate v_{rec}/V_A scaling as $\delta/L \ll 1$, and hence would not be rapid enough to be of any practical interest. Thus, we come to a conclusion that Petschek's mechanism (or a variation thereof) is necessary for sufficiently fast large-scale reconnection.

Recently, however, several numerical and analytical studies (e.g., [6–16]) and laboratory experiments [17] have shown that, in resistive MHD with a uniform (and, by inference, Spitzer) resistivity, Petschek's mechanism does not work; the slow Sweet-Parker scaling applies instead. In other words, in the collisional regime, when classical resistive MHD applies, one does not get Petschek's fast reconnection.

It is then natural to ask now whether fast reconnection is possible in a collisionless plasma where resistive MHD is not valid. The answer now appears to be "yes." First, in space and solar physics fast collisionless reconnection events have been observed for a long time. More recently, it has also been seen in laboratory experiments [17,18]. In addition, several theoretical and numerical studies have recently indicated that fast Petschek-like reconnection does indeed take place in the collisionless regime. Moreover, there may even be two physically distinct mechanisms for fast collisionless reconnection: (i) Hall effect (e.g., [9,12,19–24]) and (ii) spatially localized anomalous resistivity (e.g., [7,8,13–16,25–27]). At present, it is still not clear which of them operates under what conditions and how they interact with each other. However, both of these mechanisms seem to work and both seem to involve an enhancement due to a Petschek-like configuration. Thus, I believe it is safe to say that a Petschek-enhanced fast reconnection does indeed happen in the collisionless regime.

To summarize, there are two regimes of magnetic reconnection: the slow Sweet-Parker reconnection in resistive MHD with classical collisional resistivity and the fast Petschek-like collisionless reconnection.

Now, how can one quantify the transition between these two regimes? First, consider the case with a relatively weak (or zero) guide field, $B_z \lesssim B_0$, where B_0 is the reconnecting field component. Then, the condition for fast collisionless reconnection can be formulated (e.g., [9,15,18,24,27,28]) roughly as

$$\delta_{SP} < d_i. \tag{1}$$

(Since the discussion in this paper is very approximate, I will consistently ignore all numerical factors of order 1.)

Expressing resistivity that enters via δ_{SP} in terms of the Coulomb-collision electron mean-free path $\lambda_{e,\text{mfp}}$, one gets [18]:

$$\frac{\delta_{SP}}{d_i} \sim \left(\frac{L}{\lambda_{e,\text{mfp}}}\right)^{1/2}\left(\beta\frac{m_e}{m_i}\right)^{1/4}, \tag{2}$$

where β is the ratio of the plasma thermal pressure $(2n_eT_e)$ at the center of the layer to the reconnecting magnetic field pressure $(B_0^2/8\pi)$ outside of the layer. The condition of force balance across the layer (in the absence of a strong guide field) requires $\beta \simeq 1$, where we neglected the contribution due to the upstream gas pressure. Then, the above fast collisionless reconnection condition becomes

$$L < L_c \equiv \sqrt{m_i/m_e}\lambda_{e,\text{mfp}} \simeq 40\lambda_{e,\text{mfp}}. \tag{3}$$

Note that, by construction, the mean-free path that enters here is due to classical Coulomb collisions. It is given by $\lambda_{e,\text{mfp}} \simeq 7\times10^7$ cm $n_{10}^{-1}T_7^2$, where we took $\log\Lambda \simeq 20$ and where n_{10} and T_7 are the central layer density n_e and temperature T_e in units of 10^{10} cm^{-3} and 10^7 K, respectively. Substituting this into Eq. (3), we get

$$L < L_c(n, T) \simeq 3\times10^9 \text{ cm } n_{10}^{-1}T_7^2. \tag{4}$$

The strong T_e dependence tells us that knowing the temperature is crucial. Note that n_e and T_e that enter here are those at the center of a Sweet-Parker reconnection layer and are not known *a priori*. Therefore, we would like to cast condition (4) in an alternative form that would involve only the ambient plasma parameters, such as the far-upstream values of the plasma density, temperature, and magnetic field. Now, the cross-layer pressure balance $(2n_eT_e = B_0^2/8\pi)$, valid in the zero-guide-field case, provides us with one relationship between the central n_e and T_e, and so, by itself, it is not sufficient. Indeed, this condition only tells us that the thermal pressure at the center of the layer needs to be raised to a certain level to balance the outside magnetic pressure, but it does not tell us whether this is achieved by increasing the density or the temperature. In order to break this degeneracy, we need to consider also the equation of energy conservation. The logic of our model dictates that this analysis be done in the collisional Sweet-Parker regime. At the minimum, this analysis should include Ohmic heating and heat advection and an estimate of various possible energy-loss mechanisms, such as radiation and electron thermal conduction. In particular, it can be shown [28] that (i) on the time that a fluid element spends inside the layer (the Alfvén transit time $\tau_A = L/V_A$), Ohmic heating converts to heat just enough magnetic energy to raise T_e to the level required by the pressure balance, without the need to increase the density substantially, (ii) for the solar coronal conditions, radiative losses are negligible on the Alfvén time scale, (iii) the energy losses due to parallel electron thermal conduction can be neglected provided that the layer is collisional, in the sense of Eq. (3), and (iv) the energy losses due to the perpendicular electron thermal conduction are only marginally important at best. All these findings lead us to the conclusion that, once the condition $L > 40\lambda_{\text{mfp}}$ is satisfied and hence the system is in the collisional Sweet-Parker regime, the energy equation can be regarded basically as a balance between Ohmic heating and advection. As a result, the plasma density in the center of the layer should remain roughly comparable to the ambient coronal density, whereas the temperature should increase dramatically. In fact, Ohmic heating is enough to raise the temperature up to the "equipartition" level that depends only on the ambient density and upstream magnetic field B_0, and is insensitive to the ambient coronal temperature:

$$T_e \sim T_e^{\text{eq}} \equiv \frac{B_0^2/8\pi}{2k_Bn_e} \simeq 1.4\times10^7 \text{ K } B_{1.5}^2 n_{10}^{-1}, \tag{5}$$

where $B_{1.5} \equiv B_0/(30\text{ G})$. This estimate applies only as long as the resulting value is much higher than the ambient temperature (typically of the order of 2×10^6 K for the solar corona), which means that the ambient plasma β with respect to the reconnecting magnetic field should be $\lesssim 1$. Using this estimate, the collisionless reconnection condition can finally be written as [28,29]

$$L < L_c(n, B_0) \sim 6\times10^9 \text{ cm } n_{10}^{-3}B_{1.5}^4. \tag{6}$$

In this form, the fast-reconnection condition is particularly useful because it involves only the global quantities characterizing a given reconnecting system: its global length L, the reconnecting component of the magnetic field B_0, and the ambient plasma density n_e.

Next, let us consider the strong guide field case, $B_z \gg B_0$, which is in fact more relevant to the problem of solar coronal heating. Although some of the arguments and results presented above have to be modified, conceptually, they remain similar. In particular, the relevant collisionless reconnection scale becomes the ion-acoustic gyroradius ρ_s calculated with the total magnetic field $B_{\text{tot}} \simeq B_z$ (e.g., [30]). Correspondingly, the collisionless reconnection condition becomes [31]

$$\delta_{SP} < \rho_s \sim d_i\beta_e^{1/2}\frac{B_0}{B_z}, \tag{7}$$

where, again, β_e is based on the central n_e and T_e and on the upstream reconnecting field component B_0. Once again, all the quantities entering Eq. (7) are to be estimated in the collisional Sweet-Parker regime. To do this, first note that, in the strong guide field case one can no longer use the cross-layer pressure balance to deduce $\beta_e \sim 1$; this is because a relatively slight compression of the guide field can always ensure the pressure balance. Moreover, a strong guide field effectively makes the plasma incompressible, so that the central n_e is equal to the ambient value. But this still leaves us with the task of evaluating the central electron temperature that is needed to determine the Spitzer resistivity. It turns out, however, that the above energy-balance arguments for a collisional Sweet-Parker layer still apply, at least qualitatively [28]. Therefore, the equipartition estimate for the central temperature, given by Eq. (5), should still approximately hold; in particular, one should still have $\beta_e \sim 1$. Then, one can repeat the procedure outlined above and derive the following approximate condition for the transition to fast collisionless reconnection in the strong guide case [28]:

$$L < L_c(n_e, B_0, B_z) = \sqrt{\frac{m_i}{m_e}}\lambda_{e,\mathrm{mfp}}\left(\frac{B_0}{B_z}\right)^2 \simeq 6\times 10^9\ \mathrm{cm}\ n_{10}^{-3}B_{1.5}^4\left(\frac{B_0}{B_z}\right)^2. \quad (8)$$

Thus, the main effect of a strong guide field is to reduce the critical global length L_c by a factor $(B_z/B_0)^2 \gg 1$; that is, the collisionless reconnection condition becomes harder to satisfy. Also, it is interesting to note that, for fixed values of n_e and B_z, L_c becomes very sensitive to the reconnecting field component: $L_c \sim B_0^6$.

Let us now discuss the implications of these findings for the solar corona. As long as flux emergence and braiding of coronal loops by photospheric footpoint motions continue to generate current sheets in the corona, magnetic dissipation in these current sheets results in intermittent heating [2,32]. Typical dimensions and field strengths of these current sheets are basically determined by the emerging magnetic structures and by the footpoint motions. The main focus of this Letter, then, is on what sets the typical level of the coronal plasma density and how it relates to the intermittent nature of energy release in the corona. My main point is that coronal heating should be viewed as a self-regulating process that keeps the corona marginally collisionless in the sense of Eqs. (1)–(8) (see [28,29]).

As a first example of how this works, consider a coronal current sheet with some given fixed L, B_0, and B_z. Resolving (8) with respect to n_e, we can define a critical density n_c below which reconnection switches from the slow collisional to the fast collisionless regime:

$$n_c(B_z \gg B_0) \sim 2\times 10^{10}\ \mathrm{cm}^{-3}\ L_9^{-1/3}B_{1.5}^{4/3}\left(\frac{B_0}{B_z}\right)^{2/3}. \quad (9)$$

Notice that this value is comparable to the typical densities observed in the active solar corona. I argue that this is not just a coincidence. Indeed, if at some initial time the ambient density n_e is higher than $n_c(L, B_0, B_z)$, then the current layer is collisional and reconnection is in the slow mode. Energy dissipation then is weak; hence, the plasma gradually cools through radiation and precipitates to the surface. The density around the given current sheet drops and, at some point, becomes lower than n_c. Then, the system switches to the fast collisionless regime and the rate of magnetic dissipation jumps. Next, there is an important positive feedback between coronal energy release and the density. A large part of the released energy is conducted to the surface, where it is deposited in a dense plasma. This leads to chromospheric evaporation along the post-reconnected magnetic loops, filling them with a dense and hot plasma. The density rises and may now exceed n_c. This will shut off any further reconnection (and hence heating) involving these loops until they again cool down, which occurs on a longer, radiative time scale. Thus we see that, although highly intermittent and inhomogeneous, the corona is working to keep itself roughly at the critical density given by Eq. (9). In this sense, coronal heating is self-regulating [29].

As a second example, consider a situation in which the initial density is relatively low, so that radiative cooling rate is much slower than the footpoint twisting rate. Then one can regard the density as constant between reconnection events and focus instead on the slow evolution of the reconnecting magnetic field, caused by the motion of the footpoints (similar to Refs. [2,33]). Let us consider, for example, a flux tube, anchored on the solar surface at both ends, with a fixed strong axial (guide) field B_z and a fixed volume V. The tube is composed of smaller flux fibrils that are being wrapped around each other by the photospheric motions. Over time, this wrapping leads to the formation and strengthening of current layers. For simplicity, let us represent this process by a linear growth of the reconnecting field component of a single current sheet of a fixed length L: $B_0(t) = \gamma t B_z$. Here, γ parametrizes the rate of twisting. Let us now try to follow the evolution of this system. At first, B_0 increases steadily in time, while the density stays constant. This continues until B_0 reaches a critical value that depends on n_e according to collisionless reconnection condition (for fixed L and B_z):

$$B_c(n_e, L, B_z) \sim 30\ \mathrm{G}\ L_9^{1/6}n_{10}^{1/2}B_{z,2}^{1/3}, \quad (10)$$

where $B_{z,2} \equiv B_z/(100\ \mathrm{G})$. We assume here that B_0 always stays well below B_z. As soon as this critical value is reached, the system switches to the fast-reconnection regime and magnetic energy $B_0^2/8\pi$ is rapidly dissipated. Importantly, part of this energy is not radiated promptly but is transported by parallel thermal conduction to the solar surface. This causes an evaporation episode adding new plasma to the flux tube under consideration. The amount of plasma added is roughly proportional to the energy released in a given event, $B_0^2/8\pi \sim B_c^2(n_e)/8\pi$, which, according to (10), is in turn proportional to the density in the tube just before reconnection: $\delta n_e \sim B_c^2(n_e) \sim n_e$.

Now let us see what happens on a still longer (several hours) time scale. As a consequence of the first fast-reconnection event, the current sheet is promptly destroyed and B_0 drops back to nearly zero. The field-line twisting, however, still continues, and so the process described above repeats. This time, however, the plasma density in the tube is higher, and hence the critical magnetic field B_c

is larger and takes longer time to reach. In particular, taking the twisting rate γ to be constant, the time between subsequent reconnection events scales as $\delta t = \gamma^{-1} B_c(n_e)/B_z \sim n_e^{1/2}$. Therefore, as long as the relative increase in density at each step is small, the long-term ($t \gg \delta t$) evolution can be effectively described by the differential equation

$$\frac{dn_e}{dt} \simeq \frac{\delta n_e(n_e)}{\delta t(n_e)} \sim \sqrt{n_e}, \tag{11}$$

and so $n_e(t) \sim t^2$. Correspondingly, the emission measure of the tube increases as t^4.

This growth will continue until one of the following two effects intervenes. First, as the density builds up, the critical value of B_0 may become so large (a sizable fraction of B_z) that the equilibrium shape of the entire loop will be affected. The loop may then undergo the kink instability and become sigmoidal, which, with further twisting, may result in a large-scale eruption with a catastrophic energy release (a large flare).

Alternatively, it may happen that the density just builds up gradually to a level large enough for radiative cooling between two subsequent reconnection events to become important. Indeed, as the density increases, the radiative emission measure increases as n_e^2 and the time δt between reconnection events as $n_e^{1/2}$ (see above). Ignoring for simplicity coronal temperature variations, the amount of thermal energy lost between two reconnection events scales as $n_e^{5/2}$, whereas the amount of thermal energy gained after each reconnection event is just proportional to $B_c^2(n_e) \sim n_e$. At some point, the two will inevitably become comparable. Correspondingly, the amount of plasma drained due to the gradual radiative cooling will become equal to that pumped back up into the corona by each chromospheric evaporation episode. Then, on some long time scale (but still only as long as γ, L, and B_z remain constant), the evolution of the system can be represented by repeated cycles that include fast-reconnection events, followed by chromospheric evaporation episodes, followed by relatively long (~1 h) periods during which the free magnetic energy builds up and the plasma gradually cools down.

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