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Nanoflare heating model for collisionless solar corona

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Abstract. The problem of coronal heating remains one of the greatest unresolved problems in space science. Magnetic reconnection plays a significant role in heating the solar corona. When two oppositely directed magnetic fields come closer to form a current sheet, the current density of the plasma increases due to which magnetic reconnection and conversion of magnetic energy into thermal energy takes place. The present paper deals with a model for reconnection occurring in the solar corona under steady state in collisionless regime. The model predicts that reconnection time in the solar corona varies inversely with the cube of magnetic field and varies directly with the Lindquist number. Our analysis shows that reconnections are occurring within a time interval of 600 s in the solar corona, producing nanoflares in the energy range 10^{21}–10^{23} erg/s which matches with Yohkoh X-ray observations.

Keywords. Magnetic reconnection in solar corona; nanoflare heating; fast magnetic reconnection; solar corona.

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1. Introduction

How the nature of the processes that heat the solar corona, maintain it at this high temperature and accelerate the solar wind is a great mystery still unsolved fully. The high temperature of the corona requires a permanent heating mechanism. From many observations and recent simulations it is clear that magnetic fields play an important role in the heating process [1,2]. The corona contains a large assembly of high-temperature elemental magnetic strands in the coronal loops which are rooted at the photosphere. Magnetic reconnection, which is the process that enables us to convert magnetic energy into heat and kinetic energy of accelerated particles, is believed to be a fundamental process in various energetic phenomena in space and astrophysical systems such as solar flares [3,4]. Although the idea of magnetic reconnection for explaining the energy release in solar flares and coronal heating had been proposed many decades ago [5,6] it was after Yohkoh observations [7] that the role of magnetic reconnection occurring during energy release in the solar corona was identified as a strong possibility. Recent observations from SOHO, TRACE and RHESSI have been giving further evidence for magnetic reconnection in the solar corona.

The high-resolution coronal imager (Hi-C) launched on a sounding rocket on 11 July 2012 took images of the 1,500,000-K corona with a resolution of 0.2 arcsec (roughly 150 km). The data obtained by Cirtain *et al* [8] from Hi-C observations show evidence of small bundle of loops that are wrapped or braided about each other along their length, reconnection and consequent heating in the lower solar corona. They also observed a small flare at the intersection of the converging strands and emergence of plasma from the reconnection region.

In the present paper, we have considered the coronal plasma as collisionless due to its high temperature and propose a new analytic model for magnetic reconnection in the solar corona which successfully explains the observed heating. Our result shows that reconnections occurring in the solar corona at a time interval of 600 s cause the solar flares. From each reconnection event, the energy of the order of 10^{21} erg is carried as thermal energy outwards from the coronal plasma in each second. This energy is sufficient for active coronal heating in a localized region.

2. Proposed model

2.1 *Sweet–Parker reconnection*

The Sweet–Parker model uses a steady-state scaling analysis to determine how reconnection parameters vary with system plasma parameters. The Sweet–Praker reconnection rate in the resistive MHD regime can be calculated as shown below.

When two magnetic fields of opposite polarity come very close, they reconnect each other and a current sheet is formed. This current sheet of finite region is called diffusion region. The energy of the stressed field is released by reconnection. As more and more magnetic field lines are moving towards diffusion region, thermal pressure increases, and plasma flows from the diffusion region with speeds in the order of Alfven speed (figure 1).

$v_{\mathrm{outflow}} = v_{\mathrm{A}} = B/\sqrt{\mu_0 m_{\mathrm{i}} n}$, assuming steady-state condition and assuming the fluid to be incompressible. Here B is the ambient magnetic field in the coronal plasma, m_{i} is the mass and n is the number density of ions.

Applying continuity equation,

$$v_{\mathrm{inflow}} L = v_{\mathrm{A}} \delta_{\mathrm{sp}}, \quad (1)$$

where L is the length and δ_{sp} is the width of the diffusion region in the Sweet–Parker mode.

The electric field occurring in this region can be explained using the generalized Ohm's law,

$$E = -v_{\mathrm{e}} \times B - \frac{\nabla\psi}{n_{\mathrm{e}} e} - \frac{m_{\mathrm{e}}}{e}\left(\frac{\mathrm{d}v_{\mathrm{e}}}{\mathrm{d}t}\right) + \eta j,$$

where E and B are the electric and magnetic fields, v_{e} is the electron velocity, ψ is the pressure tensor, η is the resistivity, j is the current density, m_{e} is the mass and e is the electron charge.

When the total derivative is split, we get

$$E = -v_{\mathrm{e}} \times B - \frac{\nabla\psi}{n_{\mathrm{e}} e} - \frac{m_{\mathrm{e}}}{e}\left(\frac{\partial v_{\mathrm{e}}}{\partial t} + (v_{\mathrm{e}}\nabla) v_{\mathrm{e}}\right) + \eta j. \quad (2)$$

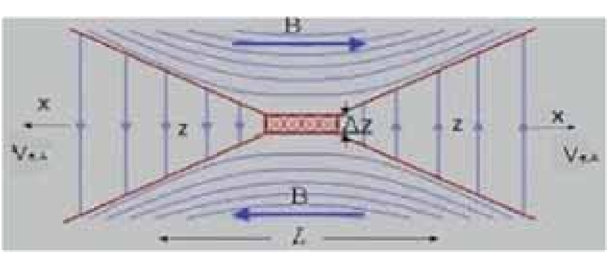

Figure 1. Basic 2D reconnection model.

In the resistive MHD regime, the advection $v_{\mathrm{e}} \times B$ and Ohmic loss (ηj) can be equated and all other terms in eq. (2) remain insignificant [9].

$$v_{\mathrm{inflow}} B = \eta j,$$

$$v_{\mathrm{inflow}} B = \frac{\eta B}{\mu_0 \delta_{\mathrm{sp}}}. \quad (3)$$

$j = (\nabla \times B)/\mu_0$ is used with $\nabla = 1/\delta_{\mathrm{sp}}$. The reconnection rate in the Sweet–Parker region is defined as

$$R_{\mathrm{sp}} = \frac{v_{\mathrm{inflow}}}{v_{\mathrm{A}}},$$

where v_{inflow} is the velocity of the reconnection inflow and v_{A} is the Alfven velocity [10]. Substituting the value of $1/\delta_{\mathrm{sp}}$ from eq. (1) and after simple calculation we get

$$R_{\mathrm{sp}} = \frac{v_{\mathrm{inflow}}}{v_{\mathrm{A}}} = \frac{1}{\sqrt{S}}, \quad (4)$$

where

$$S = \frac{\mu_0 L v_{\mathrm{A}}}{\eta}, \text{ the Lindquist number.}$$

2.2 *Invalidity of Sweet–Parker reconnection in solar corona*

For Sweet–Parker reconnection, we get the reconnection time as

$$t_{\mathrm{R}} = \frac{\tau_{\mathrm{A}}}{R_{\mathrm{sp}}} = \tau_{\mathrm{A}}\sqrt{S}$$

from eq. (4). The solar coronal parameters $B = 100$ G, $T = 10^6$ K, $n = 10^{15}$ m^{-3}, $L = 10^7$ m give Alfven velocity as 7×10^6 m/s and the corresponding Alfven time is

$$\tau_{\mathrm{A}} = \frac{L}{v_{\mathrm{A}}} = 1.42 \text{ s}.$$

For solar corona, the Lindquist number obtained is of the order of 10^{14} [11], which leads to very slow reconnection rate and is physically invalid for solar flares. We get the reconnection time t_{R} as 0.3 days for this model. Further, the Sweet–Parker model is based on resistive MHD and it cannot be applied when kinetic effects come to play. So this model fails in the case of solar flares.

2.3 *Fast collisionless reconnection model*

In this scenario, we have considered a collisionless reconnection for the solar corona. Kinetic effects play a significant role in collisionless magnetic reconnection and particle distribution functions in diffusion region are essentially non-gyrotropic. The reconnecting electric field is given by Hesse *et al* [12],

$$E_{\rm rec} = -\frac{1}{n_{\rm e}}\left(\frac{\partial P_{xy}}{\partial x} + \frac{\partial P_{zy}}{\partial z}\right), \tag{5}$$

where P_{xy} and P_{zy} are eletron pressure tensor non-gyrotropic components.

We assume $(\partial/\partial z) \gg (\partial/\partial x)$ because the width of the current sheet is much smaller than its length. Therefore, we get

$$E_{\rm rec} = -\frac{1}{n_{\rm e}}\left(\frac{P_{zy}}{\Delta z}\right), \tag{6}$$

where Δz is the width of the current sheet in the collisionless regime and also assuming $(\partial/\partial z) \approx (1/\Delta z)$.

The electron pressure tensor non-gyrotropic component P_{zy} can be expressed as [13]

$$P_{zy} = -\frac{P_{zz}}{2\Omega_{\rm e}}\frac{\partial v_{\rm ez}}{\partial z}. \tag{7}$$

Here $v_{\rm ez}$ can be considered as $v_{\rm e,\,inflow}$, the electron inflow speed because now magnetic field is advected by the electrons, and

$$\Omega_{\rm e} = \frac{eB_0}{m_{\rm e}}\frac{r_{\rm L,e}}{\Delta z},$$

where it was assumed that the possible electron bounce widths are severely constrained by the electron Larmor radius ($r_{\rm L,e}$) given by

$$r_{\rm L,e} = \frac{v_{\rm th,e}}{\omega_{\rm ce}}.$$

Here $v_{\rm th,e}$ is the electron thermal speed and $\omega_{\rm ce}$ is the electron cyclotron frequency in the guide magnetic field. If the field lines are antiparallel, the characteristic length scale is equivalent to the Larmor radius for a particle travelling in Alfven speed [14].

Therefore,

$$P_{zy} = -\frac{P_{zz}}{2\Omega_{\rm e}}\frac{\partial v_{\rm e,inflow}}{\partial z}, \tag{8}$$

where P_{zz} is the gyrotropic pressure component.

When P is the pressure in the middle of the diffusion region and P_0 is the gas pressure at large distances from the diffusion region, assuming the pressure is balanced, we can write

$$P - P_0 = \frac{B_0^2}{2\mu_0}. \tag{9}$$

The pressure difference across the diffusion region can be equated with the gyrotropic pressure component P_{zz} because outside the diffusion region, the gas pressure is smaller, while inside the pressure is higher. So fluid is ejected along the field lines reducing the built-up pressure and allowing the oppositely directed field lines to come closer. Thus, during the formation of current sheet the magnetic energy is converted into thermal energy.

Substituting the value of gyrotropic pressure component, eq. (8) becomes

$$P_{zy} = \frac{-B_0^2}{2\mu_0}\frac{m_{\rm e}}{2eB_0}\frac{\Delta z}{r_{\rm L,e}}\frac{v_{\rm e,inflow}}{\Delta z}. \tag{10}$$

In the collisionless regime, $E = -v_{\rm e} \times B$. Hence taking eq. (6) we can balance the advection $v_{\rm e} \times B$ with electron pressure tensor.

$$E_{\rm rec} = v_{\rm e,inflow}B_0 \approx -\frac{1}{n_{\rm e}}\frac{P_{zy}}{\Delta z}. \tag{11}$$

Applying continuity equation like eq. (1) we can write

$$v_{\rm e,inflow}L = v_{\rm e,A}\Delta z, \tag{12}$$

where L is the length of the current sheet and $v_{\rm e,A}$ is the electron Alfven speed given by

$$v_{\rm e,A} = \frac{B}{\sqrt{\mu_0 m_{\rm e} n}}$$

where $m_{\rm e}$ is the mass and n is the number density of electrons. The value of electron Alfven speed obtained is 2.8×10^8 m/s.

Substituting (10) in (11) and also using (12) we get

$$v_{\rm e,inflow}B_0 = -\frac{1}{n_{\rm e}}\frac{B_0^2}{2\mu_0}\frac{m_{\rm e}}{2eB_0}\frac{v_{\rm e,A}}{r_{\rm L,e}L}. \tag{13}$$

The collisionless reconnection rate is found to be

$$R_{\rm collisionless} = \frac{v_{\rm e,inflow}}{v_{\rm e,A}} \approx \frac{1}{2ne^2}\frac{m_{\rm e}}{2\mu_0 r_{\rm L,e}L}. \tag{14}$$

We can calculate the reconnection time as

$$t_{\rm R} = \frac{\tau_{\rm A}}{R_{\rm collisionless}}$$

using the Alfven time

$$\tau_{\rm A} = \frac{L}{v_{\rm e,A}}.$$

We get

$$t_R = \frac{4\mu_0}{v_{e,A}} \frac{ne^2 r_{L,e}}{m_e} L^2. \quad (15)$$

By taking the length of diffusion region in terms of Lindquist number,

$$L = \frac{\eta S}{\mu_0 v_{e,A}}$$

and substituting in eq. (15),

$$t_R = \frac{4\mu_0}{v_{e,A}^3} \frac{ne^2 r_{L,e}}{m_e} \frac{\eta^2 S^2}{\mu_0^2}. \quad (16)$$

Substituting for Alfven speed, eq. (16) becomes

$$t_R = \frac{4e^2\eta^2 S^2 n^{5/2} (m_e\mu_0)^{1/2} r_{L,e}}{B^3}. \quad (17)$$

Equation (17) is the main result of our model which gives reconnection time in terms of magnetic field, Lindquist number and Larmor radius. From eq. (17) we find that the reconnection time is inversely proportional to the cube of the magnetic field.

The rate of energy release per unit time in the current sheet is obtained by

$$\frac{dE}{dt} = \frac{B^2 V}{2\mu_0 t_R}, \quad (18)$$

where V is the volume of the current sheet given by $L^2\Delta z$. On substituting the values of eq. (17) in eq. (18) we get

$$\frac{dE}{dt} = \frac{B^5 L^2 \Delta z}{8e^2\eta^2 S^2 \mu_0^{3/2} m_e^{1/2} n^{5/2} r_{L,e}}. \quad (19)$$

From eq. (19) we found that the dissipative energy per unit time in a current sheet varies as the fifth power of the magnetic field.

2.4 *Application to solar corona*

In the collisional region, we start with a current sheet of 10^6 m thickness and 10^7 m length. The thinning of current sheet occurs due to instabilities and once a small scale is achieved, fast reconnection occurs because anomalous resistivity now sets in [15]. Taking the values of coronal plasma as mentioned above and also considering anomalous resistivity as $\eta = 0.1705$ Ω-m in solar corona [16], we find Lindquist number as $S = 10^9$ using the equation

$$L = \frac{\eta S}{\mu_0 v_{e,A}}.$$

Let us suppose that the current sheet is passing through a region of marginal collisionless condition. Then the thickness of the current sheet is $\Delta z = 10^3$ m [17] and length is of the order of 10^6 m. Using electron thermal speed and electron cyclotron frequency, the value of Larmor radius obtained is 2.2 mm in solar corona. Substituting this, the reconnection time from eq. (17) is of the order of 600 s which matches with observational results.

Numerically, from eq. (19) we get

$$\frac{dE}{dt} \approx 10^{14}\ \text{J/s} \approx 10^{21}\ \text{erg/s}$$

i.e., the energy release from the magnetic reconnection, in steady state, obtained in the range of nanoflares. It matches with Yohkoh X-ray observations by Katsukawa and Tsuneta [18].

3. Conclusions

In this paper, a new analytic model of collisionless reconnection is proposed. The reconnection time is found to vary inversely with cube of magnetic field. Further, it also depends on Lindquist number. Our model predicts a reconnection time of 600 s which matches with the observations of impulsive phase of the solar flares. The model predicts an energy dissipation of 10^{21} erg/s from the coronal plasma in each reconnection event, which is in the range of nanoflares. This energy range is also the same as observed by Katsukawa and Tsuneta where they found evidence in the Yohkoh X-ray observations of myriads of brightness fluctuations of the order of 10^{21}–10^{23} erg over times of 10^2 s, opening the way for direct quantitative observational studies of the phenomenon. Hence, our present model sufficiently explains active coronal heating in a localized region.

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