CHAPTER 2
ROLE OF MHD IN INITIATION OF CME
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2.1 Introduction

CME, the phenomena which are driven magnetically but pressure and gravity forces also play a role in destabilization of CMEs, like the solar wind which is driven by pressure and gravity forces (Parker, 1963). It is not possible to measure directly coronal magnetic fields, in detail this is the main reason for the problem of CME initiation which has not been yet solved. The CMEs are driven by the energy in the magnetic field and the most intriguing question in solar physics is how a magnetic field configuration could release energy in the solar corona as to create solar explosions such as CME and solar flare. Various models are given to explain the initiation of CME phenomenon.

2.2 The CME Precursors

The pre-CME structures are called CME precursors (or the CME progenitors) may be the unstable or metastable coronal structures which are very important to predict the occurrence of a CME and to understand CME triggering mechanism. In some CMEs open magnetic field (coronal hole) is the progenitor whereas in others the strongly twisted or sheared magnetic structures are the precursors. The closed magnetic field on the sun consists of active regions and bipolar magnetic fields straddling over quiescent filaments are often the source regions for CMEs. Since some active region filament and almost all quiescent filaments are of inverse polarity type are explained by flux rope model so flux rope system (A sheared and twisted field embedded in less sheared magnetic system) is the ideal model for the progenitor of CMEs (Leroy et al.,


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The advantages of the flux rope system as a CME precursor’s are-

1. This is fundamental model for twisted magnetic field lines, which carry electric current and magnetic free energy.

2. This system matches the three part structure of CMEs.

3. This system gracefully explains the eruption of CMEs with magnetic reconnection and without magnetic reconnection.

A flux rope is like arcade of twisted magnetic field lines, coming out from the positive polarity from photosphere, take one and more turn in the corona making magnetic dips with the inverse polarity, and then go back to negative polarity in

![Figure 2.1: A three-dimensional (3-D) representation of the coronal magnetic field drawn as solid colored lines at t = 0 hours. The flux rope is drawn with blue and red lines showing a sheared (toroidal) core surrounded by a highly twisted sheath, respectively. Orange and yellow lines show the poloidal field of the steady-state equatorial streamer belt. On the X–Z plane, the computational mesh is drawn with black lines superimposed upon a false color image the velocity magnitude (Manchester et al., 2004).](image-url)
the photosphere (Priest et al., 1989). These dips of magnetic field lines hold a thread of filament. Whenever an arcade of strongly-sheared flux tubes (called core field) rises due to either magnetic rearrangement or instability, the overlying less sheared bipolar magnetic loops (called envelop fields) is stretched up, and the magnetic separator collapse into the current sheet under core field. The magnetic reconnection of this current sheet cause the rapid eruption of core field into interplanetary space. Here are some observational precursors that precede CMEs.

1. **Helmet streamers swelling /slow rise of prominence** found by Hundhausen (1993) by SMM satellite observation.


3. **Soft X-Ray(SXR) brightening** found by Simnett & Harrison(1985)

4. **Radio noise storms (sometimes called type I radio burst)** found by Lentos et al.,1981

5. **Filament darkening and widening**

6. **Long term filament /prominence oscillations**

7. **Outward-moving blob near the edge of streamers.**

### 2.3 CME Triggering Mechanism

It is widely accepted that the required energy to power a strong CME comes from coronal magnetic field. The several triggering mechanism have been proposed either conceptually or through MHD analysis and /or on the basis of simulation results are described below:

2.3.1  **Tether-cutting or Flux Cancellation Mechanism**

2.3.2  **Shearing Motions Mechanism**

2.3.3  **Magnetic Breakout Mechanism**

2.3.4  **Emerging Flux Triggering Mechanism**
2.3.5 The Equilibrium Loss Mechanism
2.3.6 Instability and Catastrophe-related Triggering Mechanisms
2.3.7 Other Mechanisms.

2.3.1 Tether-cutting or Flux Cancellation Mechanism

Moore and LaBonte proposed this mechanism after analysis of filament eruption event on 29 July, 1973. The field lines that provide the tension are sometimes called tethers, analogous to the ground-anchored ropes that hold down a buoyant balloon. According to this mechanism a filament is supported on dips by magnetic field that is nearly aligned with magnetic inversion line (Fig. 2.2). Just around the filament, the magnetic fields are strongly sheared (field lines AB and CD). AB and CD are like tethers constraining the filament. These strongly sheared core field is overlaid by less sheared envelop magnetic arcade. As the magnetic shear increases, the negative lag of the field line AB is about anti-parallel to positive leg of the field line CD which causes the formation of a strong current sheet occurs with micro-scale instability-driven anomalous resistivity and magnetic reconnection.

![Figure 2.2: The tether-cutting triggering mechanism for CMEs. Strongly sheared core field is restrained by the overlying less-sheared envelope field (left); The reconnection between field lines AB and CD triggers the core field to rise (middle); and the rising core field stretches up the envelope field, forming a current sheet below the core field (Right) (Moore et al., 2001).](image)

This magnetic reconnection cut the line tied field lines AB, CD and forming a long field line AD (concave upward) and short loop CB (concave downward) following the
reconnection outflow, the AD expands upward and CB shrinks downwards (Wang, 2006). As the localizes reconnection goes on, the core field near AD pulls up the filament to rise further stretches the envelope magnetic field with the formation of elongated current sheet above magnetic neutral line. This is the triggering phase of whole CME eruption. Thus, reconnection below the cavity in this mechanism removes the overlying field that keeps the pre-existing flux rope in static equilibrium.

Flux cancellation model is same as tether cutting model in nature. The only difference is that flux cancellation model emphasizes a more gradual evolution of magnetic reconnection in the photosphere (Fig. 2.3) whereas tether cutting is more impulsive process occurring in low corona. Flux cancellation model was proposed by Ballegooijen and Martens (1989).

**Figure 2.3:** Flux cancellation in a strongly sheared magnetic arcade leading to the formation and levitation of a flux rope. Further cancellation leads to the eruption of the flux rope (Ballegooijen & Martens, 1989).

### 2.3.2 Shearing Motion Mechanism

Shearing motion is an effective way to increase the magnetic shear near polarity inversion line. This motion is important for building the free energy for the corona. The 2.5-dimentional MHD simulations performed by Mikic & Linker (1994) in the spherical coordinates indicate that due to the localized shearing motion the magnetic field expand outward and stretches the field lines and produce tangential discontinuity (a strong current sheet) above polarity inversion line. Once the resistivity excited in
the current sheet, a magnetic reconnection cause the release of magnetic energy with the ejection of plasma.

2.3.3 Magnetic Breakout Mechanism

This model was proposed by Aly and Sturrock (1991). The magnetic breakout model is a quadrupolar structure with two adjacent arcades, having overarching magnetic field lines over the whole system that represents the tethers. One loop arcade is continuously sheared and builds up magnetic stress until magnetic reconnection starts in the overlying X-point between the two loop arcades. Thus, the initial magnetic configuration consists of quadrupolar topology, with a null point being above the central flux system. As the central flux system experiences shearing motions, it expands upward, pressing the X-type null point to form a current layer. If gas pressure and resistivity are considered, the current sheet would undergo magnetic reconnection. The magnetic reconnection process removes the constraint of higher magnetic loops, then opens up the magnetic field in an upward direction (i.e., the “break-up” phase). In the final state the central flux system becomes fully open. As the central flux system raises, a current sheet forms underneath, whose reconnection leads to the drastic formation and eruption of a flux rope.

A further development of the 2D quadrupolar model of Aly & Sturrock (1991) is the so-called magnetic breakout model of Antiochos et al. (1999b) and Aulanier et al. (2000b), which involves the same initial quadrupolar magnetic configuration, but undergoes an asymmetric evolution with the opening up of the magnetic field on one side. In this magnetic breakout model, reconnection removes the unsheared field above the low-lying, sheared core flux near the neutral line, which then allows the field above the core flux to open up (Antiochos et al., 1999b). It circumvents the Aly–Sturrock energy limit by allowing external, disconnected magnetic flux from a neighboring sheared arcade (which is not accounted for in the “closed-topology” model of Aly and Sturrock) to assist in the opening-up process.
Figure 2.4: The evolution of the magnetic field in the breakout model, showing the reconnection above the central flux system removes the constraint over the core field (thick lines), and results in the final eruption (Antiochos et al., 1999).

2.3.4 Emerging Flux Triggering Mechanism

The process of flux emergence has been considered as a driver in the model of Heyvaerts et al. (1977). This model consists of three phases: (1) a pre CME heating phase starts when a new magnetic flux emerges beneath the filament and continuously

Figure 2.5: Flare model of Heyvaerts et al. (1977). (a) A reconnection occurs continuously between the old and new flux which results equilibrium loss due to heating of current sheet at a critical height and (b) A marginal reconnection between the existing flux and emerging flux where the current sheet reaches a new steady state.
reconnection occurs between the old and new flux that heats the current sheet; (2) the impulsive phase where the heated current sheet loses equilibrium at a critical height and turbulent electrical resistivity causes the current sheet rapidly to expand, results triggering chromospheric evaporation; and (3) the main phase where the current sheet reaches a new steady state with marginal reconnection.

A requirement of this model is the pre-existence of a stable current sheet (with very low resistivity) for periods of the order of a day or more. However, numerical simulations indicate that the current sheets reconnect almost as quickly as they are formed (Forbes & Priest, 1984; Shibata et al., 1990).

Chen & Shibata (2000) proposed an emerging flux triggering mechanism for CMEs where magnetic reconnection between the emerging flux and the pre-existing coronal field either inside or outside the filament channel.

![Figure 2.6: Schematic diagram of the emerging flux triggering mechanism for CMEs.](image_url)

(a) Emerging flux inside the filament. (b) Emerging flux outside the filament channel reconnects with the large coronal loop, which results in the expansion of the loop. The underlying flux rope then rises and a current sheet forms near the magnetic null point (Chen et al., 2008).
When the reconnection-favorable emerging flux appears inside the filament channel [panel (a), Fig.2.6], it cancels the small magnetic loops near the Polarity Inversion Line (PIL) below the flux rope. Thereby, under the pressure gradient, plasmas (initially in equilibrium) on both sides of the PIL are driven to move convergent along with the frozen-in anti-parallel magnetic field. As a result, a current sheet forms below the flux rope, and the flux rope is also triggered to move upward slightly.

When the reconnection-favorable emerging flux appears outside the filament channel [panel (b)], it reconnects with the large-scale magnetic loops that cover the flux rope. The right leg of the large-scale magnetic loop (rooted very close to the PIL) is re-connected to the right side of the emerging flux and becomes further from the PIL. The magnetic tension force pulls up the magnetic loop (the larger thick black line) to move upward, with the flux rope following immediately. As the plasma is evacuated below the flux rope, the X-type null point collapses to form a current sheet.

2.3.5 The Equilibrium Loss Mechanism

The driver in the Heyvaerts model is emerging flux, but the onset of CME instability is not quantified in terms of a magnetic field evolution. An evolutionary model that starts with a stable (force-free) magnetic field configuration and converging flows as a continuous driver has been developed by Forbes & Priest (1995) in 2D to find that how the (force-free) evolution passes a critical point where the system becomes unstable and triggers CMEs.

In two dimensional MHD models the eruption properties are well established, both analytically (Forbes & Isenberg, 1991) and numerically (Forbes, 1991) and suggest that an eruption will occur provided that the anchoring of the ends of the flux rope to the solar surface does not prevent it (Antiochos et al., 1999). The flux rope cannot escape unless the static arcade field associated with the line current is removed (Roussev et al., 2003). Magnetic reconnection in the current sheet releases most (≈ 95%) of the magnetic energy that has been built up from the initial force-free configuration. This fully analytically model yields reasonable amounts of released energies, suitable to explain flares and CMEs.
2.3.6 Instability and Catastrophe-related Triggering Mechanisms

When the coronal magnetic field subjected to the photospheric motions and flux emergence evolving in a quasi-statical way (both ideally and non-ideally), then a critical stage comes where the equilibrium is unstable (i.e., instability) or no nearby equilibrium state exists (i.e., loss of equilibrium).

2.3.6.1 Kink Instability

The MHD evolution of sheared or twisted flux tubes or coronal loops is quite complex. Due to the cylindrical geometry of coronal loops, current pinch instabilities represent an important class of MHD instabilities in the solar corona, such as the kink mode, sausage mode, or helical/torsional mode. Thus, kink instability is a class of MHD instabilities which sometimes develop in a thin plasma column carrying a strong axial current. If a kink begins to develop in such a column the magnetic forces
on the inside of the kink become larger than those on the outside, which leads to the
growth of the perturbation. The column then becomes unstable and can be displaced
into the walls of the discharge chamber, causing a disruption.
Considering modes that are periodic in axial direction $z$ or azimuth angle $\phi$, axially
symmetric perturbations $\xi(x)$ can be written as,

$$\xi(x) = \xi(r)e^{(kz+m\phi)} \quad \ldots (2.1)$$

Where $m = 0$ is called the sausage mode and $m = 1$ the kink mode. The $m = 0$ is
independent of the azimuth angle $\phi$ and thus represents a purely radial oscillation,
while $m = 1$ involves an azimuthal asymmetry (i.e., a sinusoidal oscillation in a
particular azimuthal plane $\phi$). The kink instability can occur in flux tubes that have
azimuthal magnetic fields $B_\phi$ (or associated azimuthal currents $j_\phi$) above some
critical threshold. If a kink-like displacement deforms a straight fluxtube, the
azimuthal magnetic field lines move closer together at the inner side of the kink than
at the outer side, which creates a magnetic pressure difference $\nabla (B_\phi^2/8\pi)$ towards the
direction of the lower field (which is the outer side of the kink), and thus acts as a
force in the same direction as the kink displacement, and therefore makes it grow
further. For force-free magnetic fields of uniform twist, a critical twist of $\phi_{\text{twist}} \sim
3.3\pi$ (or 1.65 turns) was found to lead to kink instability, while the critical value
ranges between $2\pi$ and $6\pi$ for other types of magnetic fields. The kink instability is
believed to be an important trigger for filament eruption, flare initiation, and CMEs.

![Figure 2.8](image)

**Figure 2.8:** The evolution of the kink instability of a twisted flux tube based on an
analytical solution (Sakurai, 1976).
2.3.6.2 Torus Instability

Fan & Gibson (2007) studied the emergence of a flux rope from the subsurface into the magnetized corona, with 3D MHD simulations. When the background magnetic field declines slowly with height, a kink instability occurs; whereas when the background magnetic field declines rapidly with height, a weakly-twisted flux tube, whose twist is below the threshold for kink instability, still erupts with little writhing like a plenary outward expansion. They interpret the latter case as the “torus instability”.

Figure 2.9: Two MHD simulations of the eruption of a twisted flux rope in the corona triggered by the onset of the torus instability and the kink instability.

2.3.6.3 Catastrophe (Magnetic Non-Equilibrium)

The studies of MHD catastrophe of coronal flux rope systems have confirmed the possibility that the magnetic energy stored in the corona is released by a global magnetic topological instability, which is essentially an ideal MHD process. When the instability takes place in a catastrophic manner, the plasma is accelerated by the Lorentz force. As a result, the magnetic energy is mainly transformed into the kinetic energy of plasma with the formation of current sheets. Alternatively, the 2D analytical solution given by Forbes & Isenberg (1991) indicates that as flux cancellation continues near the magnetic neutral line, the flux rope embedded in a bipolar field initially rises smoothly [Fig. 2.10 (a-d)] and at a critical point, the flux rope presents a catastrophic behavior [the transition from a null-point (at x = y = 0) to a current
sheet]. The energy released in this transition is trivial so it might be insufficient to power an eruption. Since the higher state contains a current sheet so reconnection would result in the rapid eruption of the filament. Therefore, the **flux cancellation-induced transition** works as an efficient triggering mechanism for CMEs.

Figure 2.10: Variation of the equilibrium state of a flux rope system as the amount of the cancelling flux ($\phi$) increases. From panel (e) to panel (f), a catastrophe takes place (Forbes & Isenberg, 1991).

Forbes & Priest (1995) found that even without flux cancellation, a flux rope system subject to the photospheric converging motion would also experience a catastrophic behavior. When half distance of the dipole decreases from 4 to 0.97, the flux rope always possesses only one equilibrium state. However, at 0.97, the flux rope has two equilibrium states, i.e., a higher energy state and a lower energy state with a current sheet below the flux rope. Such a catastrophe can serve as a nice trigger mechanism for CMEs. However, the 3D MHD numerical simulations by Amari et al. (2003a) indicate that as the converging motion is imposed at the bottom boundary, the flux rope always goes up. When series of deriving equilibrium state are subjected to changing parameters in the analytical solutions, the frozen-in effect in the ideal MHD might be violated. It is also seen from their comparison that the catastrophe existing in the analytical study is not present in the MHD simulations. It might be due to that
there is a “toroidal force” in the 3D flux rope (not existing in the 2D model), which excites the “torus instability” (Török & Kliem, 2007).

2.3.7 Other Mechanisms

Besides the above-mentioned triggering models, there are some other mechanisms that have not been investigated extensively and quantitatively.

2.3.7.1 Mass Drainage

It is generally assumed that filaments are supported by the Lorentz force against gravity. If a part of filament material drains down to the chromosphere, the filament would lose its equilibrium under the excess Lorentz force (or called magnetic buoyancy, Tandberg-Hanssen, 1974). Such a process was studied by Fan & Low (2003), Wu et al. (2004) and Zhou et al. (2006).

2.3.7.2 Sympathetic Effect

Morton waves and/or EIT waves generated by some CME events might trigger eruption (Ballester, 2006) of a remote filament. Below an erupting CME, the reconnection inflow may also induce the loss of equilibrium of a neighboring filament, as shown by Fig. 2.11 (Cheng et al., 2005). In addition, the primary CME pushes aside the background magnetic field, which can also induce the loss of equilibrium of a neighboring filament (Ding et al., 2006).

![Figure 2.11](image-url)

**Figure 2.11:** Schematic sketch showing that the reconnection inflow in one CME eruption induces the loss of equilibrium of a neighboring filament (Cheng et al., 2005).
2.3.7.3 Solar wind

The CME source region with a closed magnetic configuration is often bounded by open magnetic field, where solar wind is accelerated from being nearly static to several hundred km/s. It is possible that the CME source region might be pulled up by the solar wind with the drag force (Forbes et al., 2006) or the pressure gradient. The most important condition for the significant dragging effect of the solar wind is that the flux tube extends to a height where the plasma is around unity.

2.3.8 Standard Model for CME

The most widely accepted standard model for CMEs is the 2D magnetic reconnection model that evolved from the concepts of Carmichael (1964), Sturrock (1966), Hirayama (1974), Kopp & Pneuman (1976), called the CSHKP model. The model was called “standard model” for CME/flares (Hudson & Cliver, 2001).

In the model of Sturrock (1966), a helmet streamer configuration was assumed to exist at the beginning of a CMEs, where the tearing-mode instability (induced by foot point shearing) near the Y-type reconnection point triggers a CME, accelerating particles in a downward direction and producing shock waves and plasmoid ejection in an upward direction.

Figure 2.12: Temporal evolution of a flare according to the model of Hirayama (1974), (a) which starts from a rising prominence, (b) triggers X-point reconnection beneath an erupting prominence and (c) the draining of chromospheric evaporated, hot plasma from the flare loops and plasmoid ejection.
Hirayama (1974) explains the pre CME process as a rising prominence above a neutral line (between oppositely directed open magnetic field lines), which induces a magnetic collapse on both sides of the current sheet after eruption of the prominence. This magnetic collapse is accompanied by lateral inflow of plasma into the opposite sides of the current sheets and X-type reconnection occurs. This region is assumed to be the location of major magnetic energy dissipation, which heats the local coronal plasma and accelerates non-thermal particles. As a result of this impulsive heating, chromospheric plasma evaporates and fills the newly reconnected field lines with over dense heated plasma, which produces soft X-ray-emitting flare loops and CMEs.

Kopp & Pneuman (1976) refined this scenario further and predicted a continuous rise of the Y-type reconnection point, due to the rising prominence. As a consequence, the newly reconnected field lines beneath the X or Y-type reconnection point have an increasingly larger height and wider foot point separation.

The standard eruptive-flare model frequently referred to as CSHKP after Carmichael (1964); Sturrock (1966); Hirayama (1974); Kopp & Pneuman (1976) describes the global evolution of a solar eruption consisting of a flare and a CME. The initial driver of the CME process is a rising prominence/filament above the neutral line in a flare-prone active region. This rising filament stretches a current sheet above the neutral line, which is prone to Sweet–Parker or Petschek reconnection (Fig. 2.13). In a bipolar active region, the two polarities are connected by a magnetic arcade that overlies an initially stable magnetic flux rope, corresponding observationally to a sigmoid/filament/prominence. When the flux rope becomes unstable, it rises in the corona and the overlying field lines stretch outward. A current sheet forms within the arcade below the flux rope and above the polarity inversion line (Lin et al., 2005, 2008; Reeves et al., 2008). Eventually, magnetic reconnection begins in this current sheet. The resulting reconnected field lines wrap around the flux rope, building it up further to form the CME.
2.4 Role of MHD in Initiation of CME

The interaction between plasma and a magnetic field can be modeled according to the principles of MHD, in which the plasma can be treated as a continuous medium. The equations of MHD unify the equations of slow electromagnetism and fluid mechanics. The main branches of MHD are: equilibrium, waves, instabilities and reconnection. Among the various scenarios given for the initiation of CMEs - loss of equilibrium, instability or magnetic reconnection plays a crucial role. Thus, MHD is an essential tool to understand and quantify the energy released by magnetic field configurations in the solar corona. In the framework of MHD, the dynamo problem can be stated as the search for solutions of the MHD equations (the combination of the non-relativistic
Maxwell equations, Ohm’s law and the hydrodynamical equations) with non-
vanishing magnetic energy as time goes to infinity.

2.4.1 Ideal MHD Equations

MHD is expressed in terms of macroscopic parameters, such as density, pressure, 
temperature and flow speed of the plasma. This plasma reacts to (macroscopic) 
electric and magnetic forces as described by the Maxwell equations. However, 
particle motion in plasma can also be described by microscopic physics, called kinetic 
theory, such as in terms of the Boltzmann equation. The fluid approach of MHD can 
be derived from kinetic theory by defining appropriate statistical (average) quantities (Aschwanden, 
2006).

Many astrophysical plasmas are characterized by a set of equations that is called ideal 
MHD equations and includes the MHD continuity equation, the momentum equation, 
Maxwell’s equations, Ohm’s law and a specialized equation of state for energy 
conservation (e.g., incompressible, isothermal, or adiabatic). Thus, a full set of ideal 
MHD equations includes (for an adiabatic equation of state):

MHD continuity equation

\[ \frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0. \]  \hspace{1cm} (2.2)

Standard form of MHD equation, which corresponds to the hydrodynamic 
momentum equation, except that the Lorentz force (third term in right hand side) and 
viscosity force (fourth term in right hand side) is added.

\[ \frac{\rho \mathbf{Dv}}{\partial t} = -\nabla P - \rho g + (\mathbf{j} \times \mathbf{B}) + \mathbf{F}_{\text{visc}} \]  \hspace{1cm} (2.3)

Where \( \frac{D}{Dt} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla \) is convective time derivative. This equation is known as equation 
of motion where amperè’s current \( \mathbf{j} = (1/4\pi) (\nabla \times \mathbf{B}) \).

\[ \frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}) + \eta \nabla^2 \mathbf{B}. \]  \hspace{1cm} (2.4)

The equation of state is,
\[ \frac{D}{Dt} (P \rho^{-\gamma}) = 0. \] ... (2.5)

\[ \nabla \times E = -\frac{1}{c} \frac{\partial B}{\partial t}. \] ... (2.6)

\[ \nabla \cdot B = 0, \] ... (2.7)

and Ohms law

\[ E = -\frac{1}{c} (\nu \times B). \] ... (2.8)

Here \( \rho, \eta, \nu, B, P, n, \gamma \) and \( c \) are the total mass density, the resistivity coefficient, the velocity, the magnetic field vector, the total thermal pressure, the electron number density, the specific heat ratio \( \gamma = 5/3 \) and the velocity of light respectively.

The usage of such a set of ideal MHD equations involves a number of implicit approximations:

1. The plasma is charge-neutral, \( \vec{E} \) = 0, which yields the Maxwell equations \( \nabla \cdot E = 0 \) and \( \nabla \cdot j = 0 \).

2. The plasma has a very large magnetic Reynolds number, which yields the electric field \( E = -\frac{1}{c} (\nu \times B) \) from Ohm’s law.

3. The non-relativistic approximation, \( \nu \ll c \), which also implies that MHD time scales are much longer than electron or ion gyro periods and the displacement current \( (\partial D/\partial t) \) is neglected, which is valid when the plasma speed \( v \) is much slower than the speed of light.

4. A highly collisional plasma, implying the MHD time scales are much longer than collisional time scales.

5. The isotropic pressure, where the pressure tensor \( (P_{ij} = \int m_i \nu_i' \nu_j' f_{\alpha} d^3\nu') \) simplifies to the diagonal elements, \( P_{ij} = \delta_{ij} P_{\alpha} \), where \( P_{\alpha} \) for an ideal gas is \( P_{\alpha} = n_\alpha k_B T_{\alpha} \), where \( k_B \) is the Boltzmann constant.

6. The total pressure is the sum of the partial pressures \( P = \sum P_{\alpha} \), which yields \( P = (n_e + n_i) k_B T \approx 2n_e k_B T \), where \( n_e, n_i \) are the number density of electron and proton.

7. Adiabatic gas with energy equation of state \( P \rho^{-\gamma} = \text{Const.} \)
Approximations in such sets of ideal MHD equations are discussed by Priest, 1982; Benz, 1993; Boyd & Sanderson, 2003.

2.4.2 Resistive MHD

In the ideal MHD approximation it is assumed that the time scales of MHD processes are much longer than collisional processes, which guarantees that all species stay close to a Maxwellian distribution all the time. Plasma with a local Maxwellian distribution has zero viscosity and heat conduction, thus the viscosity term \( F_{\text{visc}} \) and heat conduction \( \nabla F_c \) does not appear in the ideal MHD approximation. This leaves the electrical resistivity as the only remaining dissipative mechanism.

The set of MHD equations (mass, momentum, and energy conservation) with finite electrical resistivity \( \sigma \) is called resistive MHD, consisting of the equations (as a function of \( \rho, v, P \) and \( B \)):

\[
\frac{\partial \rho}{\partial t} + \nabla (\rho v) = 0, \quad \text{...(2.9)}
\]

\[
\frac{\partial \rho v}{\partial t} = -\nabla P - \rho g + (j \times B), \quad \text{...(2.10)}
\]

\[
\frac{\partial B}{\partial t} = \nabla \times (v \times B) + \eta \nabla^2 B. \quad \text{...(2.11)}
\]

Equation 2.11 is known as induction equation where \( \eta = c^2 / 4\pi \sigma \)

\[
\nabla \cdot B = 0, \quad \text{...(2.12)}
\]

And Ohms law

\[
E = \frac{1}{4\pi \sigma} (\nabla \times B) - \frac{1}{c} (v \times B). \quad \text{...(2.13)}
\]

\[
\frac{\partial p}{\partial t} + v \cdot \nabla P + \gamma P (\nabla \cdot v) = (\gamma - 1)(\nabla \cdot K_\beta T), \quad \text{...(2.14)}
\]

where

\[
P = 2nK_\beta T. \quad \text{...(2.15)}
\]

Here \( \rho, \eta, v, B, P, n, \gamma, K_\beta, c \) and \( T \) are the total mass density, the resistivity coefficient, the velocity, the magnetic field vector, the total thermal pressure, the
electron number density, the ratio of specific heats ($\gamma = 5/3$), the Boltzmann constant, the velocity of light and the temperature, respectively.

In a resistive fluid the importance of resistivity is measured by the Lundquist number

$$S = \frac{v_A L_e}{\eta}.$$  (2.16)

Where $v_A = B/\sqrt{\mu_0 \rho}$ is the Alfvén velocity, $\eta = 1/(\mu_0 \sigma)$ the magnetic diffusivity and $L_e$ a typical (global) scale length.

Alternatively, one uses the magnetic Reynolds number $R_m = v L/\eta$ where the Alfvén velocity is replaced by a typical plasma velocity $v$. (In the literature the expression ‘magnetic Reynolds number’ frequently is also used for the quantity $S$). Large values of $S$ or $R_m$, which are typical for space and astrophysical plasmas, correspond to small resistive effects. In the limit of large $S$ or $R_m$, the terms involving resistivity can be neglected and resistive MHD equations reduces to the equations of ideal MHD equations.

Many of the same implicit approximations are made in resistive MHD as in ideal MHD, except for perfect conduction and adiabatic equation of state:

1. charge-neutrality
2. Non-relativistic speeds
3. Highly collisional plasma
4. Isotropic pressure.

A lot of discussion was done by Boyd & Sanderson (2003).

### 2.5 Magnetic Reconnection

In solar atmosphere normally the plasma is attached very effectively to the magnetic field but sometimes magnetic field can slip through the plasma and reconnect in the regions where the magnetic gradients area millions times stronger than the normal. The term magnetic reconnection was introduced by Dungey (1953). Magnetic reconnection is the process whereby magnetic field lines from different magnetic domains are spliced to one another, changing their patterns of connectivity with respect to the sources. It is a violation of an approximate conservation law in plasma physics, and can concentrate mechanical or magnetic energy in both space and time.
Reconnection is at the heart of many spectacular events (CMEs, solar flare) in our solar system. The usual principal effects of magnetic reconnection are:

(i) to convert some of the magnetic energy into heat by ohmic dissipation;

(ii) to accelerate plasma by converting magnetic energy into bulk kinetic energy;

(iii) it creates large electric currents and electric fields, shock waves and current filamentation, all of which may involve to accelerate fast particles;

(iv) it change the global connections of the field lines and so affect the paths of fast particles and heat, which are directed mainly along the magnetic field.

2.5.1 Two-dimensional Reconnection

The simplest geometry in which reconnection may be described has two spatial dimensions, requiring the presence of an ignorable coordinate in three-dimensional physical space. In this section Cartesian coordinates X,Y and Z are used and it is assumed that the physical quantities are independent of Z. We will first consider steady-state and then introduce time dependence.

2.5.1.1 Steady-state Reconnection

The basic configuration of two-dimensional steady-state reconnection is shown in Fig. 2.14. All field quantities are independent of time. Also, the magnetic field B and the plasma velocity v are assumed to lie in the X-Y plane, while for the electric field a non-vanishing Z-component is admitted. The magnetic field vanishes at the origin (neutral point); viewed three dimensionally a neutral line (line on which B = 0) extends along the Z-axis.

When a new magnetic flux system is pushed towards a pre-existing old magnetic flux system, a new dynamic boundary is formed where the magnetic field can be directed in opposite directions at both sides of the boundary. The magnetic field has then necessarily to drop to zero at the boundary to allow for a continuous change from a positive to a negative magnetic field strength. Thus the balance between the magnetic and thermal pressure across the neutral boundary layer,

\[
[p_1 + \left( \frac{B_1^2}{4\pi} \right)] = p_{nl} = [p_0 + \left( \frac{B_0^2}{4\pi} \right)]
\]

\[\text{...(2.17)}\]
yields a higher thermal pressure ($P_{\text{nl}}$) in the neutral layer (where $B = 0$) than on both sides with finite field strengths $B_i$ and $B_0$. Analogously, the subscripts ‘i’ and ‘o’ refer to the center inflow and outflow points on the boundary of the diffusion region (Fig. 2.14) and the subscript ‘nl’ is used for quantities on the neutral line.

![Diagram of magnetic reconnection process]

**Figure 2.14:** Basic 2D model of a magnetic reconnection process, driven by two oppositely directed inflows (in X-direction), which collide in the diffusion region and create oppositely directed outflows (in Y-direction). The central zone with a plasma-$\beta$ parameter of $\beta > 1$ is called the diffusion region (grey box) (Schindler & Hornig, 2001).

The process of bringing two oppositely directed magnetic flux systems together will always have a finite area of first contact, which limits the extent of the neutral boundary layer and channels outflows to both sides, where the lateral inflows (driven by external forces) will create outflows along the neutral line in an equilibrium situation. The plasma-$\beta$ parameter [$\beta = P_{\text{th}}/(B_i^2/8\pi)$] becomes larger than unity in the central region (because $B_i \to 0$), so that the plasma can flow across the magnetic field lines, which is called the diffusion region, and is channeled into the outflow regions along the neutral boundary. Outside the diffusion region the plasma-$\beta$ again drops below unity and the magnetic flux is frozen-in. The end result is a thin diffusion region with width $2l$ and length $2L$ (Fig. 2.14). The whole process can evolve into
steady-state equilibrium with continuous inflows and outflows, driven by external forces. The Lorentz force creates an electric field $E$ in a direction perpendicular to the 2D-plane of the flow and a current density $J_{nl}$ in the neutral layer is associated with the electric field $E$ according to Ohm’s law,

$$E_0 = \frac{1}{c} v_i B_i = \frac{1}{c} v_0 B_0.$$  \hspace{0.5cm} \ldots (2.18)$$

which is termed as the current sheet for the diffusion region. The finite resistivity $\sigma$ requires, a treatment in the framework of resistive MHD, although the processes outside the diffusion region can be approximated using the ideal MHD equations.

For steady-state, compressible flows ($\nabla \cdot \mathbf{v} \neq 0$), it was found that the outflows roughly have Alfvén speeds,

$$v_A = v_0 = B/\sqrt{(4\pi\rho)}$$ \hspace{0.5cm} \ldots (2.19)$$

and that the outflow speed $v_0$ relates to the inflow speed $v_i$ reciprocally to the cross sections $2I$ and $2L$ (according to the continuity equation),

$$v_i \rho_i l = v_0 \rho_0 l.$$ \hspace{0.5cm} \ldots (2.20)$$

Using condition of incompressibility:

$$v_i l = v_0 l,$$ \hspace{0.5cm} \ldots (2.21)$$

and that the reconnection rate $M$, defined as the Mach number ratio of the external inflow speed $v_i$ to the (Alfvén) outflow speed $v_A$, (with the approximation $B_i \approx B_0$, $v_i \approx v_0$),

$$M = \frac{v_i}{v_A} = \frac{1}{\sqrt{S}}.$$ \hspace{0.5cm} \ldots (2.22)$$

$S$ is the Lundquist number (or magnetic Reynolds number) is defined by $S = (v_A l)/\eta$, where the diffusion region has length scales $2I$ and $2L$. The reconnection is a two parameter process, for instance described by $M$ and $S$. The parameter $M$ is usually called reconnection rate. It measures the velocity with which the plasma enters the region of consideration (normalized by the local Alfvén velocity). The reconnection rate in two dimensions is measured by the electric field at the reconnection site. This electric field is perpendicular to the plane and it prescribes the rate at which magnetic flux is transported from one topological domain to another (Vasyliūnas, 1975).
2.5.1.1.1 Sweet–Parker Reconnection

In the Sweet-Parker mechanism, it is assumed that the diffusion region is a thin extended structure such that we identify the sheet length (L) with the global external length-scale (L_e). Thus, magnetic Reynold number S is identified as the **global magnetic Reynolds number** \( S = (v_A L_e) / \eta \). The external region is largely homogeneous such that approximately \( B_1 = B_0 \) and \( S_1 = S_0 \). The reconnection rate is \( M = 1 / \sqrt{S} \) for instance, in the solar corona where S lies between \( 10^6 \) and \( 10^{12} \), the fields reconnect at between \( 10^{-3} \) and \( 10^{-6} \) of the Alfvén speed.

2.5.1.1.2 Petschek Reconnection

Petschek (1964) proposed a model with an increased rate of reconnection associated with a greatly reduced size of the diffusion region to a very compact area (L \( \approx l \)) that is much shorter than the Sweet–Parker current sheet (L \( \gg l \)). Since the length of current sheet is much shorter, the propagation time through the diffusion region is shorter and the reconnection process becomes faster. Most of the inflowing plasma turns around outside the small diffusion region and slow mode shocks arise where the abrupt flow speed changes from \( v_1 \) to \( v_2 \) in the outflow region (Fig. 2.15). The shock waves are the main sites where inflowing magnetic energy is converted into heat and kinetic energy.

![Diagram of Petschek Reconnection](image)

**Figure 2.15:** Petschek's model, in which the central shaded region is the diffusion region and the other two shaded regions represent plasma that is heated and accelerated by the shocks.
In the Petschek analysis the magnetic field decreases from a uniform value \(B_e\) at large distances to a value \(B_i\) at the entrance to the diffusion region given by
\[
B_i = B_e \left( 1 - 4 \frac{M_e}{1} \log \frac{B_e}{B_i} \right).
\] ... (2.23)

Petschek (1964) estimated the maximum reconnection rate \(M\) at a distance where the internal magnetic field is half of the external value i.e. the mechanism chokes itself off when \(B_i = 1/2 \, B_e\), which gives a maximum reconnection rate \(M_e^*\) of
\[
M_e^* \approx \frac{\pi}{8 \log R_{ne}}
\] ... (2.24)

Therefore, for coronal conditions, where the magnetic Reynolds number \(R_{ne}\) is very high \((10^8 \ldots 10^{12})\), the Petschek reconnection rate is \(M_0 \approx 0.01 \ldots 0.02\). Thus, the Petschek reconnection rate is about three orders of magnitude faster than the Sweet–Parker reconnection rate.

2.5.1.2 Unsteady/Bursty Reconnection

Although, steady-state reconnection has been widely used to explain many solar phenomenon it seems that in many cases magnetic reconnection occurs as a time-dependent process. Several features of steady-state reconnection are also present in typical time-dependent (two-dimensional) cases, such as a neutral line and an associated stagnation flow pattern. A qualitatively different case arises when reconnection occurs as an unstable process. Some of unsteady reconnection modes are, such as tearing instability, coalescence instability and their combined dynamics (i.e., the regime of bursty reconnection). There are also other unsteady reconnection types, such as X-type collapse (Dungey, 1953), resistive reconnection in 3D (Priest & Forbes, 2000) or collisionless reconnection (Haruki & Sakai, 2001b).

2.5.1.2.1 Tearing Mode Instability

As the name suggests, the tearing mode breaks or tears up the current sheet into a number of smaller magnetic islands and changes the magnetic topology of the system. This instability occurs when an increase in island size leads to a state of lower magnetic energy. This is the very important mode of the instability of a neutral sheet first pointed out by Dungey (1958). It requires dissipation like finite resistivity in the excitation. When the diffusion region gets too long (such as in the Sweet–Parker
model), it becomes unstable to secondary tearing (Furth et al., 1963) and produce impulsive bursty regime of reconnection (Priest & Forbes, 2000). At the current sheet itself, the plasma cannot move perpendicular to the field lines, but if the magnetic field becomes zero at the same points along the current sheet, the restraining force is much reduced and instability will set-in. The plasma will be driven towards such points in the current sheet due to non-uniformities in the field outside the sheet. An X-shaped neutral point develops and the sheet 'tears'. This could happen repeatedly along the length of the sheet forming 'islands' of magnetic field. The tearing mode

![Image of magnetic field topology]

**Figure 2.16:** The change of magnetic topology due to excitation of the tearing mode instability.

devlops a series of magnetic islands with corresponding X-type and O-type neutral lines give what is called an impulsive bursty regime of reconnection (Priest, 1986b), characterized by a more rapid energy release in a series of bursts as the islands form (Fig. 2.16). Tearing mode instability is very important in the study of magnetic reconnection processes in space plasmas.

### 2.5.1.2.2 Coalescence Instability

While tearing mode leads to filamentation of the current sheet, the resulting filaments are not stable in a dynamic environment. If two neighboring filaments approach each other and there is still non-zero resistivity, they enter another instability, the coalescence instability, which merges the two magnetic islands into a single one (Haruki & Sakai, 2001b). Coalescence instability completes the collapse in sections of the current sheet, initiated by tearing mode instability and thus releases the main part of the free energy in the current sheet (Leboef et al., 1982).
2.5.2 3D Magnetic Reconnection

When we go from two dimensions into three dimensions, we encounter many new aspects. For example, the structure of the null points is different, the global topology of the field is much more complex and there are several different types of magnetic reconnection. In two dimensions the magnetic field vanishes at a neutral point which may be either the "X" type or "O" type, but there is a much richer variety of magnetic 3D topologies, where 3D volumes with oppositely directed magnetic fields are divided by 2D separatrix surfaces, intersection lines of two separatrix surfaces form 1D separator lines and the intersections of separator lines form 3D null points.

2.5.2.1 Null Points

Magnetic nulls can be classified in terms of the eigenvalues of the tensor $\nabla B$. They have either one real and two complex conjugated eigenvalues or three real eigenvalues. For the latter case they are called type A for $(+, -)$ signs of the eigenvalues and type B for $(-, +)$. The eigenvectors of the complex conjugated eigenvalues, or of the real eigenvalues with the same sign, span a magnetic surface called the fan surface by Priest & Titov. The third eigenvector defines the spine. In the case of more than one magnetic null the fan surfaces of an A-type and B-type null intersect at a structurally stable magnetic field line called separator. The simplest magnetic null in three dimensions has magnetic field components

$$(B_x, B_y, B_z) = (x, y, -2z), \quad \ldots(2.25)$$

![Diagram](image)

**Figure 2.17:** Magnetic structure of three dimensional null points.
which satisfy $\mathbf{V} \cdot \mathbf{B} = 0$. The resulting magnetic field lines have the structure shown in Fig. 2.17, in which there are two families of magnetic field lines through the null point. A spine approaches the origin from above and below along the z-axis. Also a fan of field lines leaves the origin in a surface (the xy-plane). The fans of two nearby null points will in general intersect in a special curve called a separator, which is a field line that joins one null to the other and is a natural location for current sheet formation.

2.5.2.2 3D Reconnection with Null Points
At a magnetic null point, three different types of reconnection have been discovered. Three special types of 3D reconnections are spine reconnection, fan reconnection, and separator reconnection illustrated in Fig. 2.18.

In spine reconnection the current is concentrated along the spine (Priest & Titov, 1996). It may be driven by continuous footpoint motions across the fan of a null point. In the case of spine reconnection, a field line penetrates the fan surface, swirls around the spine, and reconnects at the opposite side of the fan surface and spine curve (Fig. 2.18).

**Figure 2.18:** Motions of (a) field lines and (b)–(d) flux surfaces in spine reconnection.

In fan reconnection the current concentrates along the fan, and it may be driven by continuous footpoint motions across the spine of a null (Priest & Titov, 1996). The result is a rapid counter-flipping of magnetic field lines above and below the fan,
which has been observed in the solar atmosphere by the TRACE satellite (Priest & Schrijver, 1999). In the case of fan reconnection, the field line merely swirls around the spine, rotates around the fan dome and reconnects at the other side [Fig. 2.19 (a) & (b)]. In separator reconnection [Fig. 2.19(c)] the current is concentrated along a separator, which is the intersection of the fans of two nulls and it is the magnetic field line that joins one null to another (Priest & Titov, 1996; Longcope & Silva, 1998). Separator reconnection occurs due to collapse of a separator.

![Figure 2.19: Fan reconnection showing motions of (a) fields lines and (b) flux surfaces. (c) Separator reconnection (spine and separator curves are marked with thick lines, fan surfaces with hatched areas, and 3D null points with black dots. The pre reconnection line is rendered in light grey and the post reconnection field line in dark grey).](image)

2.5.2.3 Reconnection without Nulls

Magnetic nulls are not the only places where magnetic reconnection may occur. Priest, Forbes & Demoulin (1995) stated that this may occur in layer-like regions where the potential mapping or mapping of foot points of field lines shows strong gradients and which are therefore called magnetic flipping layers or quasi-separatrix layers. Reconnection may also occur in the absence of nulls at so-called quasi-separatrix layers, where the gradient of the mapping from one boundary to another is very large (Priest & Demoulin, 1995; Hornig & Rastätter, 1998; Titov et al., 1999).
Demoulin et al., 1996 have demonstrated that non eruptive flares are often located along quasi-separatrix layers.

2.6 Magnetic Reconnection as an Energy Producing Mechanism

There are two means of magnetic energy release process involve in solar eruptions. In the framework of the flux rope catastrophe models for CMEs, the energy may be released either by

(i) An ideal MHD catastrophe which belongs to global topological instability of the system. In this case energy released without Ohmic heating, especially suitable for CMEs without associated flare (Hu et al., 2003b).  

(ii) Resistive magnetic reconnection across pre-existing and rapidly developing current sheet. This resistive magnetic reconnection releases stored magnetic energy responsible for solar flare associated CME events.

Here we have just tried to solve the energy released by resistive magnetic reconnection problem by taking neutral current sheet model.

2.7 Theoretical Estimation of Energy during Magnetic Reconnection in Current Sheet

Generally, the MHD equations can describe the time-dependent evolution of the coronal plasma, but for describing a magnetic field of corona we require stable equilibrium solutions. Such stable equilibrium is expected to be force-free, where the plasma density (the associated pressure P and gravity) have no importance. The equation of motion from the basic MHD equation in a resistive medium [where in static equilibrium, time dependence vanishes (d/dt=0) and flows are constant (v=constant)]

\[-\nabla P - \rho g + (j \times B) = 0. \]  

...\[2.26\]

Since horizontal pressure is balanced where we neglect the gravity

\[-\nabla P + (j \times B) = 0 \]  

...\[2.27\]

\[-\nabla P + \frac{1}{4\pi}[\nabla \times B] \times B = 0 \]  

...\[2.28\]
\[-\nabla \left[ P + \left( \frac{\beta B^2}{4\pi} \right) \right] + \frac{1}{4\pi B V B} = 0. \quad \ldots(2.29)\]

The first term of equation 2.29 represents the gradient of total pressure (thermal and magnetic pressure) while the second term represents the magnetic tension. We have considered without bent vertical current sheet where there is no magnetic tension, we can neglect second term

\[-\nabla [P + \left( \frac{\beta B^2}{4\pi} \right)] = 0. \quad \ldots(2.30)\]

There is no variation in the total pressure with respect to regions

\[[P + \left( \frac{\beta B^2}{4\pi} \right)] = \text{Constant} \quad \ldots(2.31)\]

Total pressure remains constant (equation 2.17).

### 2.7.1 Plasma Streams Along the Breath of Current Sheet (In X-direction)

Let the current sheet of length 2l, breadth 2b along Y and X axis and it is independently with Z coordinate (Tandberg-Hanssen & Emslie, 1988). The direction of the magnetic field is along Y-axis which is varying linearly in field as a function of the distance from the centre of current sheet

\[B = (B_X, B_Y, B_Z) = B_0 \left( 0, \frac{x}{L}, 0 \right) \quad \ldots(2.32)\]

Where L is the magnetic scale height.

![Magnetic Pressure Force](image)

**Figure 2.20:** A sketch representing both the opposite magnetic pressure in formation of current sheet.
\[-\nabla [P + \frac{B^2}{4\pi}] + \frac{1}{4\pi (B \cdot V)B} = 0. \quad \ldots (2.33)\]

The first term yields a non-zero contribution while the second term vanishes \((dB_x/dx=0, dB_y/dy=0, dB_z/dz=0)\), so that

\[-\nabla [P + \frac{B^2}{4\pi}] = 0 \quad \ldots (2.34)\]

\[-\nabla P = \left(\frac{B_0}{2\pi}\right)(0, \frac{X}{L}, 0) \quad \ldots (2.35)\]

Which express inward directed magnetic pressure corresponding to the X-direction on the left and anti-X-direction on the right. Thus, the magnetic pressure drives a lateral inflow of plasma into current sheet, unless the lateral magnetic pressure is balanced by thermal pressure from hotter plasma inside the current sheet. When change in a magnetic field B and current density J is high inside the current sheet and directed along Y-axis then a reconnection occurs. The maximum current density increases before magnetic reconnection.

### 2.7.2 Plasma Streams Along the Length of Current Sheet (In Y-direction)

The coronal magnetic field depends on the vertical height and it may be expressed as the inverse of the third power (Aschwanden, 2006).

\[B(h) = B_0 \left(1 + \frac{h}{h_0}\right)^{-3} \quad \ldots (2.36)\]

with a mean dipole depth \(h_0 = 75 \text{ Mm}\). The photospheric magnetic field strength at their footpoints varies in the range of \(B_0 = 20 - 200 \text{ G}\). In corona plasma \(\beta \) \((P_{\text{thermal}}/P_{\text{magnetic}}) < 1\) and in general

\[P_1(h) + \left\{\frac{B^2(h)}{4\pi}\right\} = P_0(h) + \left\{\frac{B_0^2}{4\pi}\right\} \quad \ldots (2.37)\]

Thus, it is clear that the region \(|Y| > 1\) (outside the current sheet) the total pressure is smaller, while \(|Y| < 1\) (inside the current sheet) achieves higher pressure. If we consider that the total pressure in a current sheet rooted in a high field region in the network \((B_0 \ll B_0)\)

\[P_1(h) + \left\{\frac{B^2(h)}{4\pi}\right\} = P_0(h) \quad \ldots (2.38)\]

The total internal pressure starts to dominate over the external pressure.
\[ B_i(h) = P_0(h) - P_1(h). \] ... (2.39)

To balance the thermal pressure there is decrease in interior magnetic field and as a result of it, fluid is ejected along the field lines (along Y-axis) reducing the built up pressure and allowing the oppositely directed field lines which come closer. A coronal null point \( (B = 0) \) will occur between the oppositely directed field components where they cancel each other.

**2.7.3 Magnetic Reconnection Inducing Instability in Current Sheet**

A random disturbance of a thin current sheet could occur in the form of current density caused by the lateral inflow of plasma and Lorentz force (opposes the flow)

\[ j = \frac{c}{4\pi} (\nabla \times B) \approx \frac{cB}{\mu_0} \] ... (2.40)

\[ F_L = \frac{j \times B}{c}. \] ... (2.41)

In the current sheet where length \( L \) is very small and current density \( j \) is high enough to produce magnetic reconnection.

For simplicity, we consider the current sheet is vertically symmetric. Thus, the plasma flow into the current sheet from both sides (along X-axis) with inflow speed \( (\nu_{in}) \) together with frozen in magnetic field lines spreading along the current sheet upward and downward (along the Y-axis) with outflow speed \( (\nu_A) \) after the field line reconnection. Where the lateral inflows (driven by external forces) will create outflows along the neutral line in an equilibrium situation. The plasma-\( \beta \) parameter becomes larger \( (\beta > 1) \) in the central region (because \( B \to 0 \)), so that the plasma can flow across the magnetic field lines, termed as diffusion region, and is channeled into the outflow regions along the neutral boundary. Outside the diffusion region the plasma-\( \beta \) again drops below unity \( (\beta < 1) \) and the magnetic flux is frozen-in. When driving force (sheared magnetic field) of inflow exceeds the opposing Lorentz force then tearing mode instability (a kind of resistive instability) occurs in the current sheet. Tearing mode instability plays a major role in triggering a CME.

If \( \nu_A \) is the Alfvén velocity of fluid along the Y-axis and \( \rho \) is the plasma density and the magnetic pressure is equivalent to the kinetic energy per unit volume. By simply using Bernoulli’s theorem;

82
\[ p_0 - p_1 = \left( \frac{1}{2} \right) (\rho v_A^2). \]  
\[ \text{(2.42)} \]

In the numerical simulation of Hirose et al. (2001) it is actually found that the major part of the stored magnetic energy is converted into kinetic energy carried away in term of CME, while only a minor part is left for heating of the associated arcade flare.

\[ v_A = \frac{B}{(4\pi \rho)^{1/2}}. \]  
\[ \text{(2.43)} \]

The Alfvén speed along Y-direction and its height dependence, at a particular location h plasma density \( \rho(h) \) and magnetic field \( B(h) \) are the function of \( h \). The density and magnetic field vary strongest along the horizontal direction of a magnetic field line. The plasma density \( \rho(h) \) decreases near-exponentially with height.

\[ \rho(h) \approx \rho_0 \exp \left( -\frac{h}{\lambda_T} \right) \]  
\[ \text{(2.44)} \]

where \( \rho_0 \approx 2.0 \pm 0.5 \text{ cm}^3 \) and thermal scale height \( \lambda_T = 47 \text{ Mm} \times \left( \frac{T_e}{1 \text{MK}} \right) \).

\[ v_A(h) = 2.18 \times 10^{11} \left( \frac{B(h)}{\rho(h)^{1/2}} \right)^{1/2}. \]  
\[ \text{(2.45)} \]

The variations of plasma density and Alfvén velocity with vertical height are shown in Fig. 2.21.

From the continuity equation in the steady-state,

\[ v_{rec} = v_A b \]  
\[ \text{(2.46)} \]

where \( v_{rec} \) (inflow speed) is inflow speed and \( v_A \) is outflow speed.

\[ v_{rec} = \frac{v_A b}{1} \]
\[ v_{rec} = \frac{8b}{(4\pi \rho)^{1/2}}. \]  
\[ \text{(2.47)} \]

According to maxwell equation

\[ 4\pi j = (\nabla \times B) - \left( \frac{1}{c} \frac{\partial E}{\partial t} \right). \]  
\[ \text{(2.48)} \]

Using non-relativistic approximation \( (v_A << c) \) the term \( \partial E/c \partial t \) becomes negligible in comparison to \( V \times B \) in above equation. The current density \( j \) in current sheet becomes

\[ j = \frac{e(\nabla \times B)}{4\pi} \approx \frac{eB}{4\pi \lambda}. \]  
\[ \text{(2.49)} \]

When two oppositely directed magnetic fields \( \pm B \) press together where the time of total energy dissipation considered as reconnection time and suggest that (Parker, 2005)
Figure 2.21: The variation of magnetic field, number density and Alfvén velocity with vertical height.
\[ j = \left( \frac{cB}{4\pi} \right) (4\pi\eta t_R)^{1/2} \exp \left( -\frac{b^2}{4\eta t_R} \right). \]  

...(2.50)

Using induction equation
\[ \frac{\partial B}{\partial t} = \nabla \times \left( \nu \times B \right) + \eta \nabla^2 B. \]  

...(2.51)

The equation represents that change in a magnetic field produces transportation (I term) and diffusion (II term). The ratio of I term to II term gives a Reynolds number \( R_m \)

\[ R_m = \frac{\nu \alpha L}{\eta} = \frac{4\pi\alpha \nu L}{c^2}. \]  

...(2.52)

In most solar phenomenon \( R_m \approx 1 \) its value is very high it means that the magnetic field moves with plasma where the magnetic field frozen to plasma energy.

The rate of energy release per unit time (as power) in the current sheet is obtained by,

\[ \frac{dU}{dt} = (B^2 \nu)/(8\pi t_R) \]  

...(2.53)

Where \( B \) is the magnetic field before reconnection, \( V = (2l)(2b)\delta = 4l b \delta \) the volume of current sheets, \( t_R \) is reconnection time and \( \delta \) is the thickness of current sheet.

From equation 2.49 and equation 2.50, we get

\[ \frac{L}{(4\pi\eta t_R)^{3/2}} = \exp \left( \frac{b^2}{4\eta t_R} \right). \]  

...(2.54)

But \( b \gg \eta \) therefore, \( \exp \left( \frac{b^2}{4\eta t_R} \right) \approx b^2/4\eta t_R \). Thus, above equation becomes

\[ t_R = \frac{\pi b^2}{4\eta L^2}, \]  

...(2.55)

\[ t_R = \frac{(16\pi b^4 \nu_A^2)}{4\eta c^4 \rho^2 R_m^2}. \]  

...(2.56)

From equations 2.47, 2.53 and 2.56.

\[ \frac{dU}{dt} = \eta c^4 \rho^2 R_m^2/2\pi b^3 v_{rec}. \]  

...(2.57)

We find out that the energy released from the reconnection in single sheet along the \( Y \)-direction is not height dependent. We consider \( \rho = 10^8 \text{cm}^3, \eta = 7.16 \times 10^{-17} \) \( (T_e=2\text{MK}), \nu_A = 6.2 \times 10^8 \text{cm/sec}, b = 100 \text{cm}, l = 0.6 \times 10^5 \text{cm} \) and \( R_m = 10^5 \)

\[ \frac{dU}{dt} = 3.96 \times 10^{32} \text{erg/sec} \approx 4 \times 10^{32} \text{erg/sec} \]
Thus, during the eruption process the increase in kinetic energy is about $4 \times 10^{32}$ erg/sec which comes due to the energy released by a current sheet during magnetic reconnection.

2.7.4 Results and discussion

Numerical simulations of magnetic reconnection have been performed by many groups (Ugai & Tsuda 1977; Yan et al., 1992; Karpen et al., 1996; Antiochos et al., 2002 and Narain et al., 2006). When two magnetic fields of opposite polarity are brought together a current sheet is produced. The accumulated magnetic free energy of stressed field is released by reconnection. We find that the energy released from the reconnection in single current sheet is height independent and have value about $4 \times 10^{32}$ erg/sec where the plasma is accelerated by the Lorentz force. As a result, the magnetic energy is mainly transformed into the kinetic energy of the plasma. This energy is responsible to offset the magnetic pressure of flux rope ahead of current sheet. The eruptive flux rope drags magnetic field lines outwards and a CME result from an initial force imbalance due to removal of a part of the plasma in flux rope.