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Theoretical Models of Solar Flares

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Abstract: Recent progress in the understanding of the basic physical picture of solar flares is discussed from a theoretical point of view, with emphasis on magnetohydrodynamic processes, such as magnetic reconnection. Several models of *CME (Coronal Mass Ejection) related flare* and *compact flare* models are critically reviewed. The role of the successive emergence of twisted flux tubes is stressed, not only for modeling *compact flares*, but also for understanding *CME related flares*.

1. Introduction

Solar flares are among the most energetic and enigmatic phenomena in the solar atmosphere. Large amounts of energy ($10^{29} - 10^{32}$ erg) are suddenly released in the corona, accelerating great quantities of nonthermal particles and heating coronal and chromospheric plasmas, resulting in transient brightenings throughout the electromagnetic spectrum, including vigorous Hα brightening. Although flares are very complex and include different processes in different events, they often appear nearly similar in Hα. This may be a key reason why flares are difficult to understand.

During the Skylab era, it became clear that there are at least two types of flares; the *large two-ribbon flare* and the *simple loop flare* (e.g., Priest 1981, 1982). More recently, during the SMM and HINOTORI projects, it was realized that there are many types of flares. Flares observed by SMM have been categorized into five types; thermal hard X-ray (HX), nonthermal HX, impulsive gamma, gradual gamma, and quiescent filament eruption flares (e.g., Bai and Sturrock 1989), while Hinotori observations have been categorized into three kinds of flares, namely type A (hot thermal), type B (impulsive), and type C (gradual hard) flares (e.g., Tanaka 1983, 1987; Tsuneta 1984).

It has also been discovered that some of large two-ribbon flares are often associated with coronal mass ejections, and that these flares seem to have the same origin as that of CMEs (e.g., Kundu and Woodgate 1986). On the other hand, there are many flares which are not associated with CMEs; such flares are relatively compact, and have global magnetic field configurations which do not seem to change. In this paper, we shall use the terms *CME related flares* and *compact flares* in place of "large two-ribbon flares" and "simple loop flares".

We shall below discuss a few models for *CME related flares* and *compact flares* separately, and discuss their merits and demerits (or remaining problems). Note that it is not possible to discuss all theoretical models of flares, because, as is well known (Priest 1982), *there are as many flare theories as there are flare theorists* !

2. CME related flares

a) Trigger mechanisms

It has now become clear that some large flares are preceded by the start of a CME (e.g., Harrison 1986; Fig. 1). That is, flares are not the origin of CMEs, but both flares and CMEs seem to have the same origin, which may likely be a kind of MHD instability occurring in the global magnetic configuration of the corona.

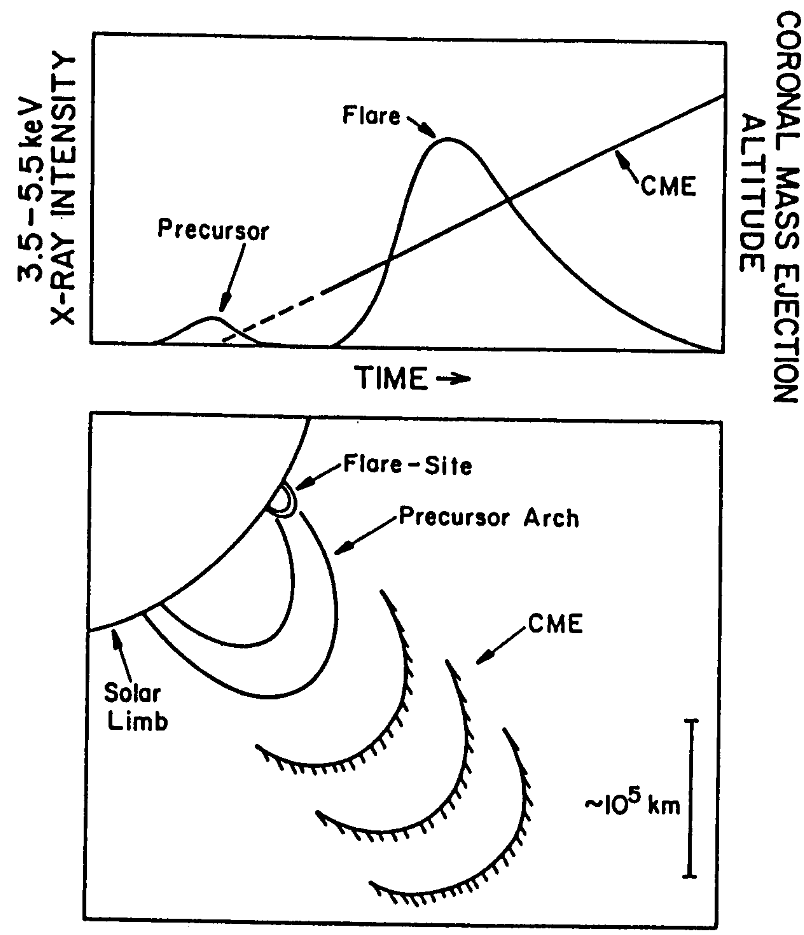

Figure 1.
The temporal relation between the start of a CME and an associated flare (Harrison 1986). Harrison writes "A coronal arch of scale-length several times 10^5 km brightens in soft X-rays (precursor). At this time a CME is launched and it appears to propagate directly from the arch. Some tens of minutes later a flare occurs in one foot of the arch."

CME related flares are often associated with filament eruptions. Such filament eruptions are believed to occur as a result of the kink instability in a twisted flux rope in the chromosphere or in the corona (Hirayama 1974, 1991; Moore 1991). Even if the filament eruption is not observed in Hα, it is possible that a twisted flux rope embedded in a hot coronal plasma becomes unstable to the kink instability, initiating a flare. Sakurai (1976) first studied the three dimensional (3D) nonlinear evolution of the kink instability, and applied the results to the dynamical motion of eruptive prominences (see also a related 3D simulation study by Zaidman and Tajima 1989).

There are three possibilities to trigger the kink instability in a twisted flux rope:

(1) Time variation of the global magnetic field configuration around the twisted rope, such as evolution induced by slow shearing motions at the footpoints of the global magnetic fields (e.g., Low 1981; Mikic *et al.* 1988; see Sakurai 1989 for reviews). (This corresponds to a change in the condition at the outer boundary of the twisted tube.)

(2) Time variation of the twisted flux tube (sheared field) itself, such as the twisting of the tube at the footpoint (e.g., Steinolfson and Tajima 1987), or reconnection leading to an unstable twisted tube (e.g., Sturrock 1989).

(3) Change in the twisted tube's lower boundary condition, such as the interaction of emerging flux with the filament from below (e.g., Rust 1972; Heyvaerts *et al.* 1977).

It is often argued that photospheric shear motions at the footpoints of the global magnetic field configuration generate the twisted flux tube or the sheared magnetic field configuration, which eventually lead to instability or the loss of equilibrium as in (1) and (2) above. It is, however, possible to interpret the observed development of sheared magnetic field configurations as the emergence of the twisted magnetic flux tube (Tanaka 1987; Kurokawa 1989). Since the energy density of the turbulent convective motion in the deep interior of the convection zone is much larger than that in the photosphere, the flux tube is much more easily twisted and sheared in the convection zone than in the photosphere (McClymont and Fisher 1989). Hence it is very possible that *twisted flux tubes formed deep inside the convection zone are the ultimate source of flares,* and that *the occurrence of flares is controlled by the emergence of such twisted flux tubes.*

Future observations should clarify this point, i.e., whether or not shear motions of footpoints of magnetic loops (arcades) actually occur. If so, there should be a pronounced global velocity field around the footpoints of the magnetic structures. .On the other hand, if the sheared configuration in the chromosphere or in the corona is a result of the emergence of already twisted flux tube from the convection zone, there shoud be an upward component in the velocity field between the two footpoints of the loop. The results of Hanaoka and Kurokawa (1989) and Kaisig *et al.* (1990) imply that the latter can be indirectly inferred from downflow velocities along the filament. Furthermore, the detailed horizontal velocity distribution around the footpoints of the magnetic loops may help to distinguish the two hypotheses, i.e., the *shear motion hypothesis* and the *emerging flux hypothesis.*

That is, if the velocity of the footpoint of a loop is very different from that of the ambient non-magnetic plasmas, the emerging flux hypothesis would be favored.

b) Flare model

One scenario for CME related flares is that following a filament eruption, previously closed field lines are opened up, and create a current sheet; i.e., a helmet-streamer type field configuration is created. If magnetic reconnection occurs successively in such a current sheet, the outward expansion of the two Hα ribbons is naturally explained (Sturrock 1968; Hirayama 1974; Kopp and Pneuman 1976).

Forbes and Priest (1983) perfomed a 2D MHD numerical simulation of magnetic reconnection occurring in a vertical current sheet line-tied at the photosphere, and found that a fast shock is created just below the downwardly directed reconnection jet. Figure 2 shows the simulation results by Ugai (1987) of loop heating by magnetic reconnection, where not only a slow shock but also a fast shock at the loop top are clearly shown. Cargil and Priest (1982) extended the model by Kopp and Pneuman (1976) to include the effect of joule heating at the slow shock front.

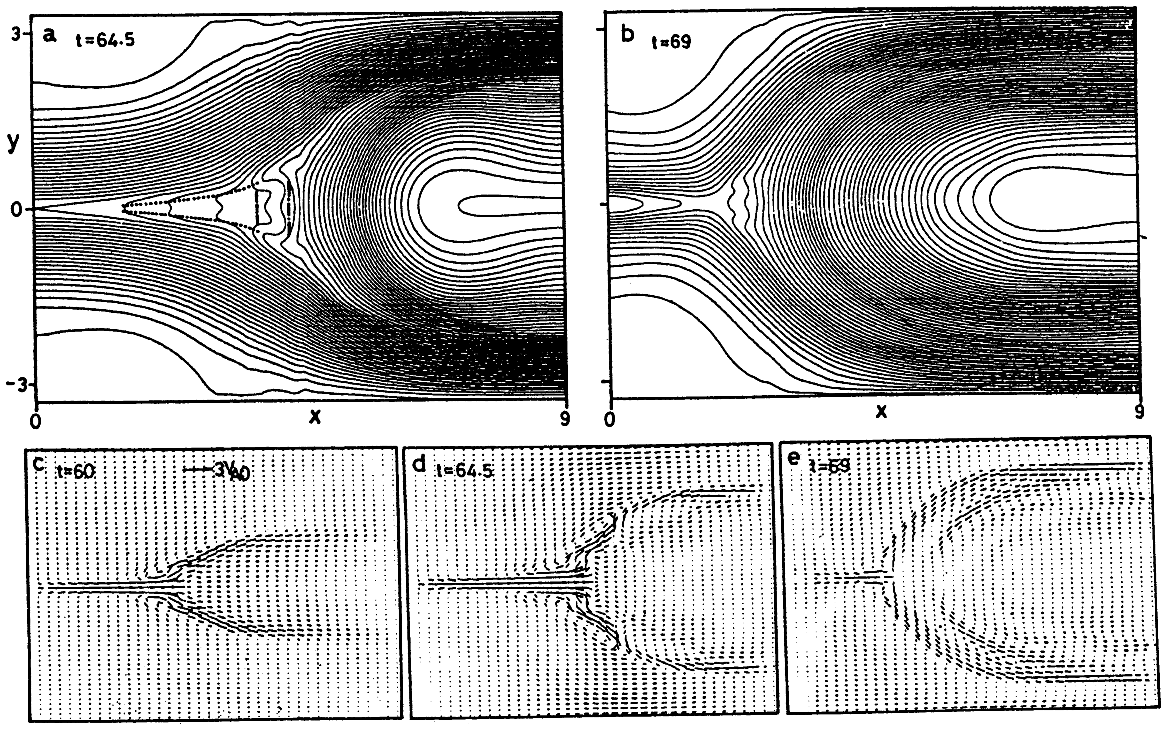

Figure 2.
Magnetic fields and plasma-flow configurations in Ugai's (1987) 2D MHD numerical simulation of loop heating by the magnetic reconnection. Both the slow shocks (dotted lines) and fast shocks (dashed) are clearly visible. The dot-dash line shows the loop front.

There are still some outstanding problems with this model: (1) There is no (obvious) external force to compress the current sheet. Modern theories on driven reconnection (e.g., Sato and Hayashi 1979; see the next subsection) have shown that an external force which compresses the current sheet is necessary to excite the fast or explosive reconnection which explains the fast rise time of flares. (2) In the model, magnetic energy stored in the twisted filament (flux tube) is not directly released in the reconnection process itself. Rather, most of the magnetic energy is converted to the kinetic energy of the filament eruption. Hence one may ask: How can such kinetic energy be converted to excite flares ? Which is essential to understand flares, the reconnection or the filament eruption ? It is to be noted here that there is some evidence from Skylab observations (Sheeley *et al.* 1975) that magnetic reconnection can occur without flares.

c) Basic theory of magnetic reconnection

Sweet (1958) and Parker (1957) have developed a simple theory of steadily driven reconnection, and found that the reconnection rate, $E = v_i B/c$, or the inflow speed, v_i, is given by $v_i = V_A {R_m}^{-1/2}$, where V_A and B are the Alfven speed and the magnetic field strength outside the diffusion region, and $R_m = V_A L/\eta$ is the magnetic Reynolds number. Here, L is the length of the current sheet (diffusion region), and η is the magnetic diffusivity. This means that the time scale of the energy release by the Sweet-Parker process is $\tau \simeq \tau_A {R_m}^{-1/2}$, where $\tau_A = L/V_A$ is the Alfven time. Since the magnetic Reynolds number is enormously large ($\sim 10^{10} - 10^{12}$) in the corona if we use the classical Spitzer (1962) conductivity, the time scale of the Sweet-Parker process is too slow to explain the impulsive phase of flares. Note that $\tau_A \sim 1 - 100$ sec for typical coronal loops, which is about the same as the time scale of the impulsive phase of flares. Even if we use anomalous resistivity ($\eta_{anomalous} \sim 10^5 - 10^6 \eta_{classical}$), the magnetic Reynolds number is still too large to explain such short time scale. This problem that the observed time scale ($\sim$ the dynamical time scale) is much smaller than the diffusion time scale is also common in fusion and magnetospheric plasma processes, and hence is one of the most challenging issues in plasma physics (Tajima 1989).

Furth *et al.* (1963) have developed tearing instability theory in an attempt to explain the abrupt disruption of magnetically confined fusion plasmas. They found, however, a time scale similar to that of the Sweet-Parker process.

Petscheck (1964) noted that considering the effect of slow mode MHD shock (or wave) on the region outside the diffusion region greatly increases the reconnection rate up to $v_i \sim (\pi/8)V_A/ln(R_m) \sim 0.01 - 0.1V_A$ (see also Sonnerup 1970), nearly independent of R_m. Although this is a very attractive idea, the size of the diffusion region is very small, so that the question arises whether the single Petcheck type reconnection controls the entire flare process or not (Kahler *et al.* 1980). Uchida and Sakurai (1977), on the other hand, studied three dimensional effects, such as the MHD interchange instability occurring in the current sheet, and proposed that the transition to a lower energy interleaved state of such an unstable current sheet may correspond to the explosive phase of flares (see also Kahler *et al.* 1980).

Petscheck's model has been confirmed by direct numerical simulations by Ugai and Tsuda (1977) and Sato and Hayashi (1979). In particular, the latter stressed the importance of externally driven reconnection in causing the sudden energy release, and suggested that any driven reconnection occurs independently of the initial R_m, even for very large R_m.

More recently, using numerical simulations of incompressible driven reconnection, Biskamp (1986) demonstrated that Petscheck's model for fast reconnection is not valid in the limit of large R_m, while Priest and Forbes (1986) developed a unified theory of steadily driven reconnection including both the Petscheck regime and the flux-pile-up regime. The latter case arises when the inflow speed exceeds that of the Petscheck regime, so that magnetic flux piles up just outside the current sheet, creating a long sheet. Forbes and Priest (1987) argued that Biskamp's result can be explained by their flux-pile-up regime, and suggested that the long current sheet appearing in the flux-pile-up regime is unstable to the tearing mode, resulting in nonsteady "impulsive bursty reconnection" (Forbes and Priest 1987; Kliem 1988). Scholer (1989) presented a somewhat different view, concurring with Biskamp (1986), that fast steady reconnection may occur if the resistivity is spatially limited so that the length of the diffusion region is sufficiently small.

Noting that nonsteady effects would be essential to understand the fast reconnection and flares (Tajima *et al.* 1982), Tajima and Sakai (1986, 1989a,b) found the explosive reconnection regime on the basis of the nonlinear simulation of the coalescence instability. According to them, if the current peaking in the magnetic island exceeds some threshold value, the reconnection of two islands proceeds in a finite time (t_0) independent of R_m; the inflow speed becomes $v_i \propto 1/(t_0 - t)$. They argued that strong driving force to compress the current sheet is essential to cause fast (explosive) reconnection. The driving force in their explosive reconnection model is the strong attractive force between two islands (current filaments).

Two current key questions in reconnection theory may be stated as follows: What are the exact conditions resulting in driven reconnection occurring independently of R_m ? And is explosive reconnection a universal phenomenon ?

3. Compact Flares

a) Emerging Flux Model

To explain compact flares, Heyvarets *et al.* (1977) developed the emerging flux model, in which the flare is produced by magnetic reconnection between an emerging loop and the overlying pre-existing coronal field.

Forbes and Priest (1984) first performed two-dimensional MHD numerical simulations of magnetic reconnection between emerging flux and a coronal field, by taking R_m to be much smaller than the actual solar value. Although their simulation results show many interesting nonlinear processes, they did not consider the effect of the gravitational acceleration. Since the main force raising the emerging flux is magnetic buoyancy (Parker 1979), the gravitational force is fundamentally important to the emerging flux model. Thus Shibata *et al.* (1989a,b,

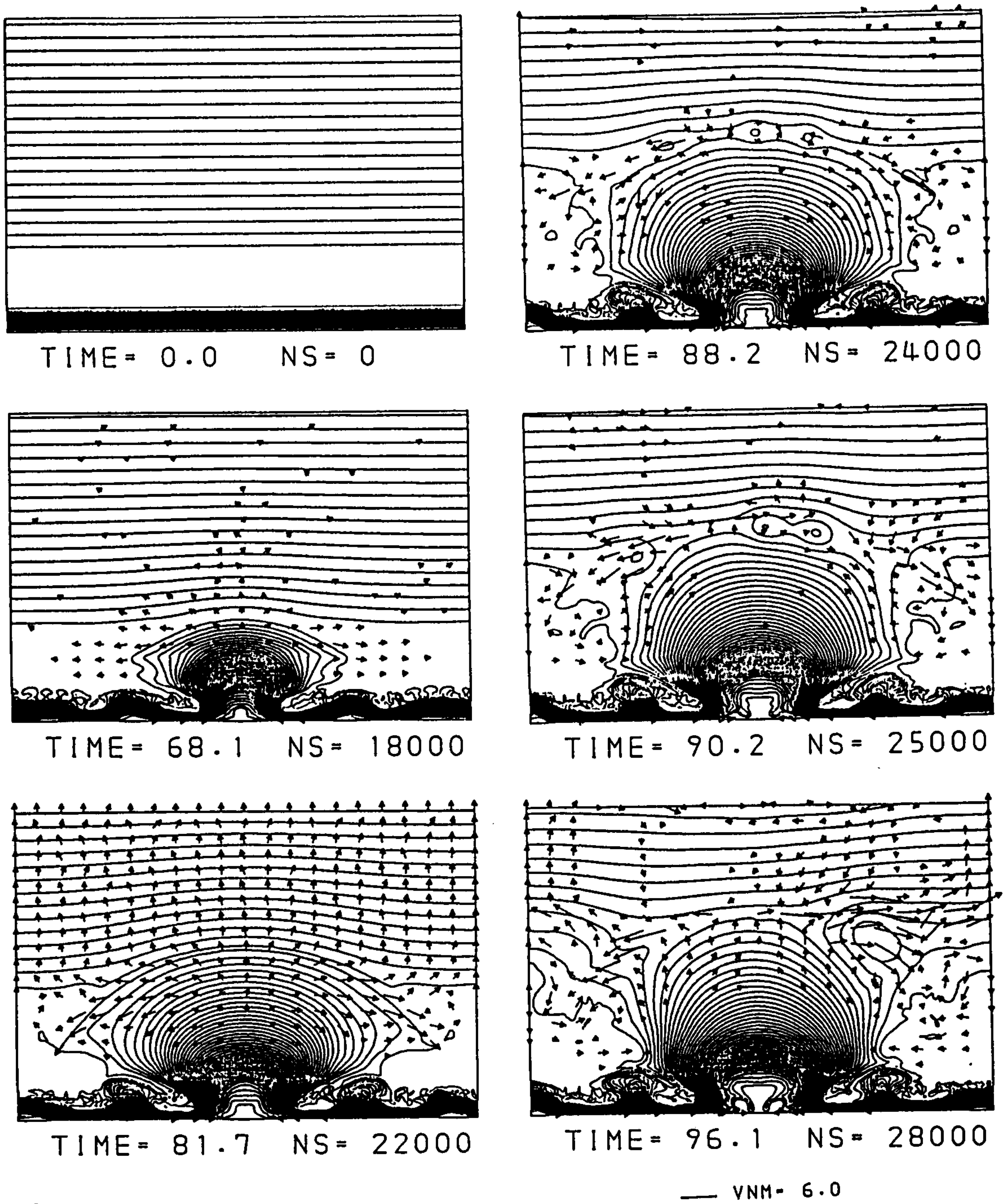

Figure 3.
Magnetic field lines and velocity vectors in a typical example of magnetic reconnection between emerging flux and an overlying coronal field (Shibata and Nozawa 1991). The times are in units of $\tau \simeq 20$ sec, and the horizontal and vertical sizes of the computing box are 16000 km and 10800 km. The scale of the velocity vector is shown below the frame of time = 96.1 in units of C_s; VNM = 6.0 indicates that the arrow with the length of this line has the velocity of $6.0C_s \simeq 60$ km s^{-1}. Note that a dense filament in the neutral sheet rises with velocity ~ 10 km s^{-1} over $t/\tau \simeq 75.1 - 81.7$. Magnetic reconnection starts after the filament gas in the sheet drains down at $t/\tau \simeq 88$. Note the formation of three magnetic islands at $t/\tau = 88.2$, their rapid coalescence at $t/\tau = 90.2$, and their subsequent jetting along the neutral sheet with supermagnetosonic speed. Hence fast shocks are created along both edges of the neutral sheet.

1990a,b; Nozawa *et al.* 1990) have constructed more realistic models of emerging flux incorporating the gravitational acceleration (see Fig. 3). Their models explain many observed features of emerging flux, such as the rise velocity ($10 - 15$ km s^{-1}) of the arch filament and the downflow speed ($30 - 50$ km s^{-1}) along the filament (Bruzek 1969; Chou and Zirin 1988), and the small rise velocity (~ 1 km s^{-1}) of the photospheric emerging flux (e.g., Kawaguchi and Kitai 1976; Zwaan 1987). (See Shibata *et al.* 1991 for a review of their recent studies on emerging flux.)

Figure 3 (Shibata and Nozawa 1991) shows a typical example of numerical simulation results of reconnection between a realistic emerging flux and an overlying coronal field, assuming small $R_m (\sim 1000)$. In this model, the resistivity is assumed to be a function of the current density and mass density, simulating an anomalous resistivity. The results show that: (1) the reconnection starts after the most of mass in the filament has fallen, (2) several magnetic islands are created in the current sheet (i.e., impulsively bursty reconnection), (3) these islands dynamically coalesce with each other via the coalescence instability (Tajima and Sakai 1989a, b), (4) the islands and neutral sheet plasmas are accelerated along the sheet up to about the Alfven speed just outside the sheet, which exceeds the local magnetosonic speed in the sheet and hence the fast shocks are created at both edges of the current sheet. Heating by the fast shocks, as well as in the current sheet, may account for the X-ray bright points associated with emerging flux (Golub *et al.* 1977). The remaining problems in these emerging flux models are the same as those of the basic reconnection theory (§2.c). It is also important to extend these 2D models to 3D models, and in particular it would be very interesting to study the emergence of the twisted flux tubes and their interaction with the overlying chromospheric and coronal magnetic fields.

The emerging flux model of Uchida and Sakurai (1977) is free from the basic difficulty of large R_m. They suggested that the very short time scale ($\sim$ dynamical time scale) of the impulsive phase of flares may be explained by the dynamical transition to a lower energy interleaved state (induced by the 3D MHD interchange instability as discussed in §2.c) of the current sheet between the emerging flux and overlying coronal field (see also Sakurai and Uchida 1977). Since the collapse is the ideal MHD process, there is no difficulty arising from large R_m. They also suggest that continued current dissipation in the interleaved book-page strucuture can provide an explanation of the later decay phase. In order to see whether these processes work well or not, 3D MHD simulations are necessary.

b) Sheared Loop and Loop Coalescence Models

Spicer (1977) has presented the sheared loop model to explain simple loop (compact) flares. He suggested that the nonlinear mode coupling and the multiple tearing modes significantly enhance the reconnection rate compared with the single tearing mode instability. Although these processes might play a fundamental role in flares, there are still some problems associated with the reconnection process, as discussed above.

Gold and Hoyle (1960) first considered the interacting (coalescence) loop model of flares. More recently, Tajima *et al.* (1982, 1987), Sakai and Tajima (1986) have presented a more refined loop coalescence model based on the concept of explosive reconnection as discussed in §2.c. They succeeded in explaining the very fast rise time of the impulsive phase of flares, and the rapid amplitude oscillations (Figs. 4) found in hard X-rays, gamma rays, and microwaves (Nakajima *et al.* 1983; Sakai and Ohsawa 1987; Tajima 1989). In addition to the basic problems discussed in §2.c, another remaining question is how the explosive coalescence evolves in three dimensional situations (see Sakai 1991).

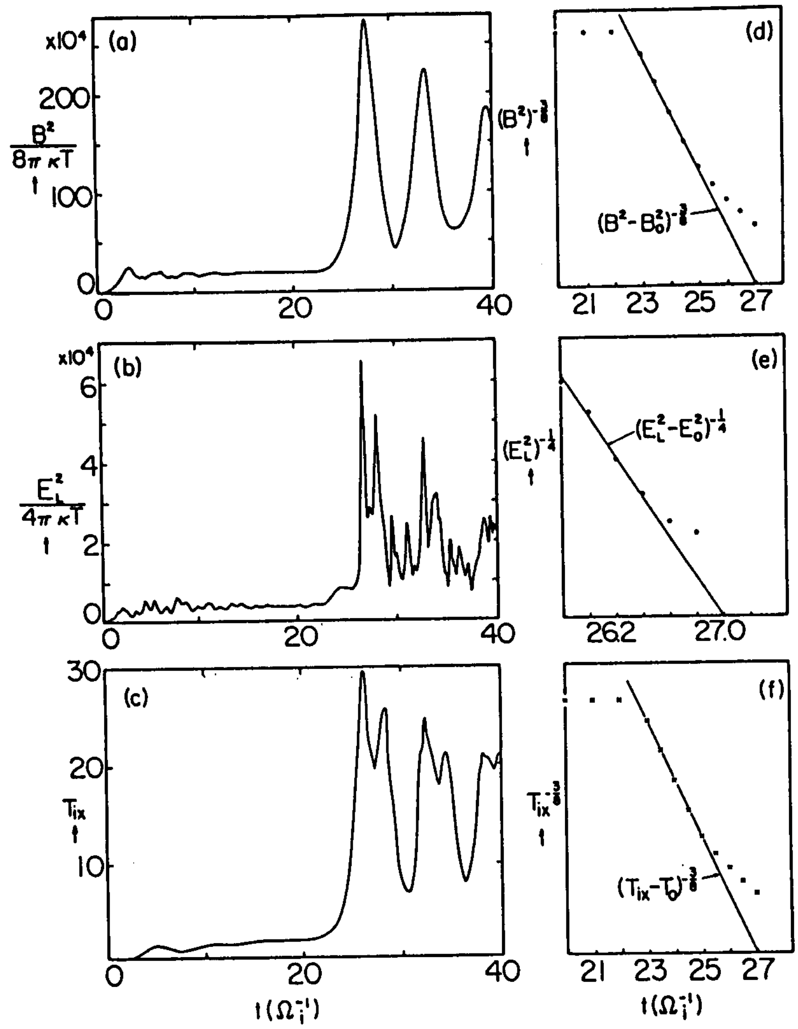

Figure 4.
Current loop coalescence model of flares by Tajima *et al.* (1987). These figures show the explosive increase of field energies and temperature during the coalescence of two magnetic islands, based on the electromagnetic particle simulations. Note that the magnetic energy, $\sim B^2$, the electrostatic energy, $\sim E^2$, and the temperature, T, diverge as $(t_0 - t)^{-8/3}, (t_0 - t)^{-4}, (t_0 - t)^{8/3}$, respectively. Note also the vigorous, large amplitude oscillations of these quantities just after the explosive phase.

c) Unwinding Magnetic Twist Jet and Loop Flare Model

Applying the mechanism of the acceleration of cosmic jets (Uchida and Shibata 1985) to solar jets, Shibata and Uchida (1986a) proposed a magnetodynamic mechanism for the acceleration of solar jets, such as surges and EUV jets (Brueckner and Bartoe 1983). They call this mechanism the ***sweeping-magnetic-twist mechanism***, or the ***sweeping-pinch mechanism***, because the acceleration is due to the $\mathbf{J} \times \mathbf{B}$ force in an unwinding (propagating or sweeping) magnetic twist (i.e., nonlinear torsional Alfven wave) and the pinching occurs in association with the propagation of the nonlinear magnetic twist (Fig. 5). This model explains very well the rotating eruption of an untwisting filament observed by Kurokawa *et al.* (1988). The origin of the magnetic twist is attributed to processes occuring deep in the convection zone, where the plasma beta is very high so that the flux tube is easily twisted by turbulent convective motion. (It is interesting to note that in the model of cosmic jets ejected from the accretion disks (Uchida and Shibata 1985; Shibata and Uchida 1986b), the magnetic twist is created by the rotation of the accretion disks.)

Uchida and Shibata (1988) then extended this mechanism to loop flares; if two magnetic-twist-jets are launched separately from the footpoints of the loop, and if the sense of the magnetic twists is opposite each other – not unlikely if we consider the emergence of the non-uniformly twisted flux tube – a very hot region appears as a result of strong shock formation when the two twists collide at the loop top. Gradual heating continues because of the dynamical relaxation of the magnetic twist which successively propagates into the region near the top of the loop (see also Uchida and Shibata 1990). An interesting point in this scenario is that it does not require any reconnection process, and thus this model is *intrinsically* independent of observed R_m. Another merit of this model is that it can explain the observed preflare upward motion along loops (Tanaka 1987).

4. Concluding Remarks

We have reviewed several models for CME related flares and compact flares, and also discussed related physical processes. In conclusion, we would like to mention the following points:

(1) The best model of the *CME related flares* may be the generalized Sturrock-Hirayama-Kopp-Pneuman (SHKP) model. It is important, however, to note that even in the framework of the SHKP model, ***the basic global process may be controlled by the successive emergence of twisted flux tubes*** as discussed by Tanaka (1987) and Kurokawa (1989). This scenario is energetically favorable to the model where the magnetic field is twisted by photospheric convective motion, because the energy density is much larger in the deep convection zone than in the photosphere. The second point regarding the SHKP model is the question of how sudden energy release occurs as a result of filament eruption. One possibility may be that multiple current loops (including the filament itself) around the filament interact with each

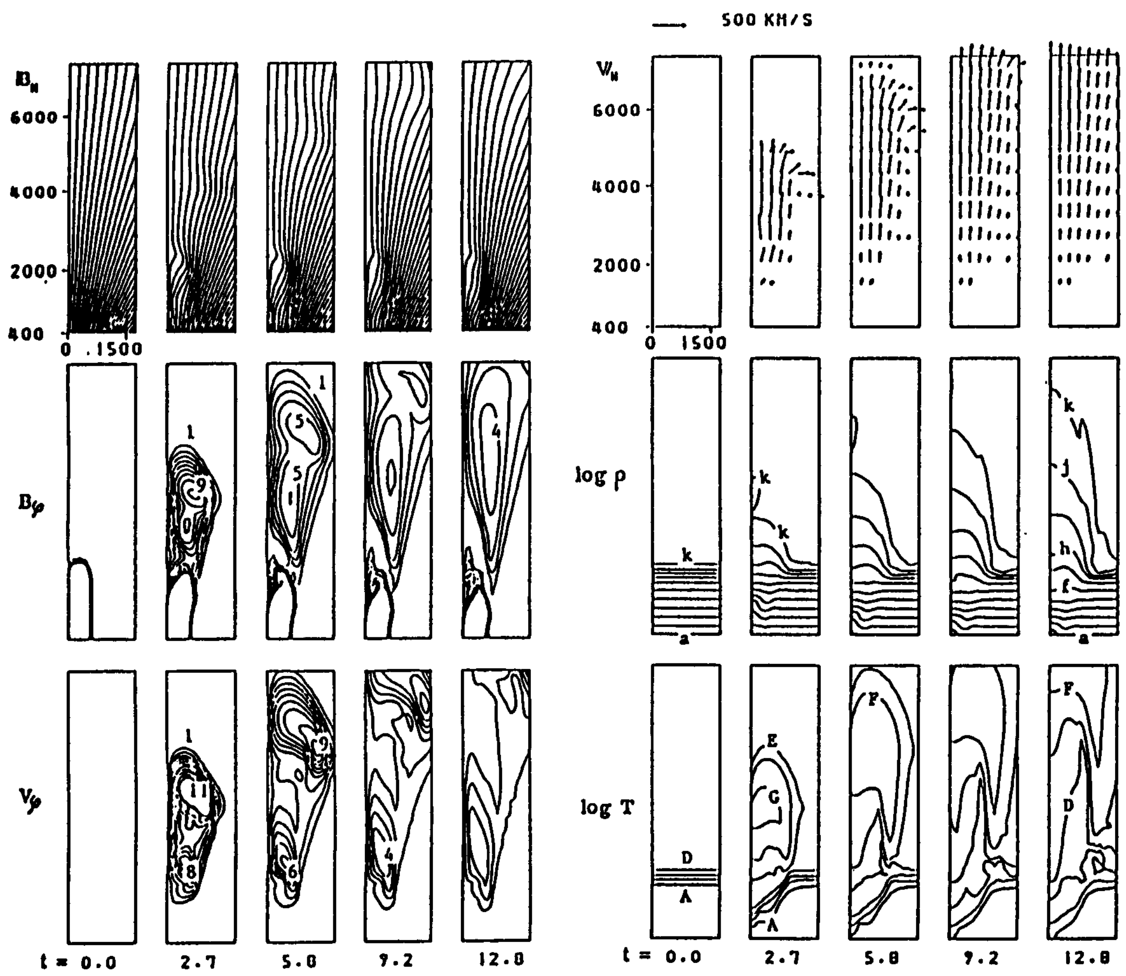

Figure 5.
The unwinding magnetic-twist-jet model by Shibata and Uchida (1986a). These figures show the results of a 2.5D axisymmetric MHD numerical simulations of the dynamical relaxation of a nonlinear magnetic twist in the upper chromosphere; from the top to the bottom in the left column, the poloidal field lines ($\mathbf{B}_{||}$), the toroidal field (B_φ) and the azimuthal (rotational) velocity (V_φ) contours; in the right column, the velocity vectors ($\mathbf{V}_{||}$), the density ($\log \rho$) and the temperature ($\log T$) contours in a logarithmic scale. The horizontal and vertical sizes of the computing box are 1600 km and 7000 km, respectively. Times are in units of seconds, and the maximum velocity of the jet is about 400 km s^{-1}. The scale of the velocity vector is shown above the frame of t = 0 of $\mathbf{V}_{||}$. The contour level step width for $\log \rho$ and $\log T$ is 0.5. Note that the jet spins about the z-axis with a rotation velocity of $\sim 60-200$ km s^{-1}. Note also that a hot region ($\sim 5-10 \times 10^6$ K, denoted by the large letter F and G) appears and propagates just ahead of the dense jet.

other as a result of the kink instability, exciting driven (explosive) reconnection (Tajima and Sakai 1989a,b).

(2) As for a model of energetic compact flares, we would like to suggest the generalized emerging flux model, in which we consider the *emergence of the twisted flux tube*. This is because a twisted tube has more free energy than a non-twisted tube, and therefore is capable of potentially producing larger flares. This model includes the previous emerging flux models (Heyvaerts *et al.* 1977; Forbes and Priest 1984; Shibata and Nozawa 1991) as a less-energetic version of compact flares, and also includes both the loop coalescence model (Sakai and Tajima 1986; Tajima *et al.* 1987) and the Uchida-Shibata (1988) loop flare model (the unwinding magnetic twist jet model).

High resolution soft X-ray imaging observations by the Solar-A project will clarify the above scenarios, and address some of the outstanding questions. At the same time, more realistic, 3D MHD simulations are highly desired to establish or reject these scenarios, and/or to resolve the questions from the theoretical side.

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