CORONAL LOOPS: CURRENT-BASED HEATING PROCESSES

P. BEAUFUMÉ
Smithsonian Astrophysical Observatory; and C.E.N. Cadarache, D.R.F.C., bat. 513, St. Paul-lez-Durance, F13108 France

B. COPPI
Massachusetts Institute of Technology, Physics Department, Cambridge, MA 02139

AND

L. GOLUB
Harvard-Smithsonian Center for Astrophysics, 60 Garden Street, Cambridge, MA 02138

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ABSTRACT

A set of new experimental observations on the structure of the solar corona (Golub et al.) is used as input to the formulation of a theoretical model of magnetic field-related coronal heating processes. Field-aligned currents are assumed to be induced along coronal loops in thin current sheaths given that the rate of current density diffusion is relatively low. The excitation of instabilities involving magnetic reconnection is proposed to occur and convert the energy associated with the current-related magnetic field directly into particle energy. The estimated heating rate in this phase exceeds the energetic requirements of the loop and the heating process is envisioned as proceeding via short bursts corresponding to an intermittent disruption of the current sheath configuration.

Because of the relatively low transverse thermal conduction, only a small fraction of the loop volume is heated to a much higher temperature than the average value. This is consistent with the experimental observations of low filling factors of hot plasmas in coronal loops (published by Martens et al.) Thus the proposed model involves a repeated sequence of dynamic events taking into account the observed loop topology, the differential emission measure distribution in the $10^6 - 10^7$ K range, the energy balance requirements in the loop, and the probable duty cycles involved in the heating process.

Subject headings: MHD — Sun: corona — Sun: magnetic fields

1. INTRODUCTION

The primary feature of the solar corona is its plasma temperature in excess of $10^6$ K, which thus radiates predominantly in the X-ray regime. Early studies used the assumption of homogeneity to derive the average properties of the solar atmosphere. However, progress in X-ray imaging techniques revealed that the corona is fundamentally inhomogeneous; for instance, Vaiana, Krieger, & Timothy (1973) identified “six classes of coronal structures observable in the X-ray photographs.”

It was initially suggested that high temperature in the corona was maintained by dissipation of various waves originating in the convection zone, but Rosner et al. (1978a) and Golub et al. (1980) suggested that coronal heating is strongly correlated with a coupling of the magnetic structures to convective motions at the surface of the Sun. More recently, the arguments against the possibility of coronal heating by purely acoustic waves and showing the central role played by the magnetic field have been summarized in Ulmschneider (1991).

Parker (1990) emphasizes that once the question of the heat input is solved, the formation and sustainment of the corona follow in a simple way. Parker shows that the heat input raises the temperature of the gas and lets it expand upward; thermal conduction then heats the upper chromosphere and evaporation fills the so-formed corona. Equilibrium with the energy losses is finally reached, the density being controlled by the radiative losses ($\propto n^2$) and the temperature being limited by the conductive losses along the magnetic field lines ($\propto T^{7/2}$). However, the observed transverse scale sizes and local contrast factors of coronal loops are not easily explained.

Given the fact that theories of the solar corona may involve spatial scales below the available resolution as well as components of the magnetic field which are difficult to observe, assessing how well-founded a given model may be is a delicate matter. A large number of models, most of them static, have been proposed to explain the heating of the solar corona. It is beyond the scope of this paper to review the theoretical state of the art in coronal heating modeling (see Narain & Ulmschneider 1990 for a recent survey; also Kuperus, Ionson, & Spicer 1981).

It is now generally accepted that the configuration of coronal magnetic structures plays an important role in transforming the energy at the photospheric level into thermal energy in the corona. However, the details of this picture are not yet firmly established. Two types of motion can generate energy in the photosphere. Periodic motions of flux tubes produce MHD waves which propagate upward and may dissipate in the corona. Also, relatively slow motions of flux tubes may be considered as responsible for d.c. field-aligned electric currents induced in magnetic loops, which can lead to their heating. Thus, the theoretical work can be divided into two branches: wave theories and release of magnetic stresses via Joule heating in magnetic reconnection (see Hollweg 1990 for a review). The presence of large-amplitude Alfvén waves in the solar wind has been known for a long time, and waves seem to be necessary to explain the plasma heating that takes place in open magnetic regions. However, it is widely believed that electric currents are an important source of coronal heating in closed magnetic structure regions. This paper will focus on the latter possibility.
Two directions have been explored in this domain:

Dissipation of currents.—Tucker (1973) showed that the current should be distributed in thin layers, and Rosner et al. (1978a) presented a model based on an anomalously high electrical resistivity generated by microinstabilities.

Generation of currents.—Parker (1972), van Ballegooijen (1986), and a number of authors (see Hollweg 1990) have developed models based on magnetic reconnection processes in configurations where the net current in the loop is nearly zero.

The purpose of this paper is to discuss possible current-based heating mechanisms, taking into account the additional observational data which are now available. In the past few years high spatial resolution X-ray observations have revealed thin loops and brightness variability which rules out the representation of these loops by static models (Sheeley & Golub 1979; Haisch et al. 1988; Raymond & Foukal 1982) and instead suggests the importance of impulsive events in the heating process (see Parker 1988). Also the three components of the photospheric magnetic field are now available by means of “vector-magnetograms.” Extrapolations of such data show the presence of high electric currents flowing along the loops in preferred directions and give estimates of the current densities and locations. Next, spectroscopic measurements have revealed the existence of very hot (~10^7 K) gas occupying a small fraction of the coronal loop volume, suggesting that the heating mechanism is very localized. The observations from the Normal Incidence X-ray Telescope (NIXT) (Golub et al. 1990) have shown that coronal loops are much thinner than was previously believed; this fact is very significant for the quantitative aspects of the theory.

In parallel with this observational progress, the study in the laboratory of magnetically confined, collisionless plasmas is very active and includes both observations and modeling for a wide range of physical phenomena. We suggest that some of these processes may be at work in solar coronal loops, and we are basing our discussion of possible coronal conditions on an extrapolation from well-established laboratory situations which have been studied for years.

Rosner, Tucker, & Vaiana (1978b) took the view that loops are the corona’s fundamental units and that each X-ray loop could be viewed as an independent “miniatmosphere.” On the basis of the recent advances in solar observations and in plasma theory, we propose a model for the heating of coronal loops whose first step is the transfer of energy from the driver (the photospheric layer) to the receiver (the corona) through the magnetic field lines. Because of the “frozen-in” effect, the footpoints of the magnetic field lines are moved around by photospheric fluid motions. We assume an electrical field along (the photospheric layer) to the receiver (the corona) through the magnetic field lines. Because of the “frozen-in” effect, the footpoints of the magnetic field lines are moved around by photospheric fluid motions.

Figure 1 (Plate 6) shows a recent high-resolution image from the NIXT rocket (see Golub et al. 1990). Such on-disk observations of the corona have made it clear that magnetic loops are a widespread feature of the corona, and in this context should be viewed as the fundamental building block of the X-ray corona (Vaiana & Rosner 1978). The coronal portion of an AR is formed by an arcade of loops of different sizes and heights, with footpoints anchored in regions of opposite magnetic polarity, the small compact loops having in general a higher
Fig. 1.—X-ray image of the solar corona obtained from a NASA sounding rocket on 1991 July 11. The X-ray corona is seen to consist of numerous long thin loops of hot plasma ($\approx 3 \times 10^6$ K) confined by the solar magnetic field.

Beaufumé, Coppi, & Golub (see 393, 397)
pressure and a higher magnetic field strength. Activity occurs within a wide range of parameters. In order to provide representative values, we shall use an artificial division of AR into long-lived versus short-lived features (Webb 1981); we shall also use values for a "typical" X-ray-bright point (XBP) even though these features are known to have a range of sizes. Little & Krieger (1977) give the following values for long-lived features (10–40 days): rise time: 1–5 days; temperature: $T \approx 2.5 \times 10^6$ K; and emission measure: $EM_{\text{max}} \sim 10^{46}$ cm$^{-3}$ for $T > 10^6$ K; for short-lived features (1–5 days): rise time: several hours; temperature: $T \sim 1.8 \times 10^6$ K; and emission measure: $EM_{\text{max}} \sim 5 \times 10^{46}$ cm$^{-3}$; and for XBP (<1 day): temperature: $T \sim 1.6 \times 10^6$ K; emission measure: $EM_{\text{max}} \sim 2 \times 10^{46}$ cm$^{-3}$.

### 2.2. Loop Parameters

Recent missions such as Skylab or SMM permitted the study of a large number of nonflaring ARs and the determination of their fundamental parameters (see Orrall 1981). Extensive observations show that there is not a single set of parameters for coronal loops but a rather wide range of values, in particular as far as the geometry and the brightness of the loops are concerned.

#### 2.2.1. Geometry

The recent pictures of the corona given by the NIXT telescope in the soft X-ray range offer very high spatial resolution (1") and reveal that coronal loops are thinner than was previously thought. These new results are available from NIXT, a sounding rocket payload utilizing multilayer coating for enhanced soft X-ray reflectivity (Golub et al. 1990); an example of the new data is shown in Figure 1. After examination of the new imaging data, we take as the basis for our calculations an ideal cylindrical loop with a ratio between the half-length $L$ (from a footpoint to the top) and the minor radius $a$ (the radius of the loop) varying from $a/L = 0.04$ for compact AR loops down to 0.01 for large-scale structures (see Fig. 2). This inverse aspect ratio is substantially lower than previous estimates, but is taken directly from the new imaging data. Coronal loop dimension measured from a number of X-ray pictures lie in the following ranges:

$$2 \times 10^{10} \leq L_{\text{[cm]}} \leq 2 \times 10^{10}$$

$$5 \times 10^{7} \leq a_{\text{[cm]}} \leq 4 \times 10^{8}$$

#### 2.2.2. Thermal Characteristics

The average temperature along the line of sight can be derived from the flux ratio of two different coronal emission lines (Vaiana et al. 1973). A number of observations and theoretical considerations allow us to establish some characteristic thermal properties for the coronal loops:

1. Along the magnetic lines, loops are mainly isothermal because of the high coefficient of thermal conduction in that direction, with sharp transition zone gradients at the ends (see Webb 1981).
2. In the radial direction, the thermal conduction is limited by the cross-field coefficient, which is much lower than the longitudinal coefficient.

The coronal loops are seen to consist of two groups (see Vaiana & Rosner 1978):

1. The brighter hot loops ($T > 2 \times 10^6$ K) are low-lying (up to $2 \times 10^8$ cm) and have a duration of several days, at least as regards the overall configuration. Their footpoints are in the area of strong or changing magnetic field, with a correlation between brighter loops and enhanced chromospheric emission. In the hot loops, hydrostatic equilibrium prevails, the loop height is smaller than the pressure scale height, and there is no clear evidence for mass flows inside the loop.
2. The fainter cool loops ($T < 10^6$ K) are longer, thinner, less stable, less numerous, and shorter-lived. There is some evidence for plasma motions in cool loops by Doppler shift measurements on SMM (Solar Maximum Mission) (see Kopp et al. 1985), and Foukal (1978) has indicated that flows must be invoked in order to explain the great height to which these loops sometimes extend.

No obvious spatial relationship appears between the two sets of loops (Cheng 1980; Webb 1981). It is to be noted that within a factor of 2 or 3, the solar corona outside of flares is essentially at a uniform temperature around $2-3 \times 10^6$ K (Golub et al. 1980; Cheng 1980), except for small amounts of very hot plasma. In active regions Pye et al. (1978) showed that...
of the EM is in this range, with relatively little cooler or hotter material. In the following, we will therefore mainly consider the hot loops as the most representative of the corona, while noting that heating of cooler loops and of the weaker quiet corona also needs to be explained.

2.2.3. Electron Density and Hot Plasma Filling Factor

The X-ray luminosity for a volume of plasma at given temperature $T$ is typically written in a form which identifies a density-independent factor $P(T)$. The total X-ray intensity is

$$L_x = \int P(T)n_e^2 \, dV,$$

(1)

where $P(T)$ is a function of the temperature $T$, $n_e$ is the electron density, and $V$ is the total volume of the emitting plasma. The emission measure (EM) is defined as $E_{\text{M}} = \int n_e^2 \, dV$. Note that if the emission is dominated by plasma occupying only a small fraction $\phi$ of the volume, we have $\langle n_e^2 \rangle V = n_e^2 \phi V$, with $\phi$ the filling factor of the gas at temperature $T$, so that $V\phi$ is the volume of the plasma considered.

Other diagnostics based on temperature and density-sensitive lines allow the determination of the density $\langle n_e \rangle$ alone, provided the temperature is known, by taking the ratio between two line intensities. Comparison between the two methods yields the filling factor $\phi$ of the plasma considered. Such measurements have been made for flaring regions in the last few years (De Jager et al. 1983) but are unfortunately more scarce for nonflaring ARs. Actually, in the flaring case the problem is of a different kind: for flares the filling factor refers to the fraction of the volume occupied by magnetic loops in the flaring region, whereas in our case we are talking of the fraction of very hot gas embedded in a single coronal loop.

The body of available evidence indicates that, unlike the temperature, the coronal density shows large variations between different loops and also within a single loop at different times. This can be understood with the Rosner et al. (1978b) scaling law

$$T \propto (pL)^{1/3}.$$

(2)

For a given $L$, $p \propto nkT \propto T^3$, so that $n_e \propto T^2$. This is a thermodynamic relation, independent of the heating mechanism. In ARs, the range is roughly $10^9 < n_{e(\text{em})} < 10^{10}$ (see Cheng 1980 and Golub et al. 1980) for a typical coronal plasma around $2-3 \times 10^6$ K. Cheng (1980) claims that the only difference between a loop and its background lies in the density and suggests that the heating mechanism is the same at any place in the corona. He proposes that it is the local enhancement of the density which makes a loop brighter than the background plasma, since the EM (and therefore also the $L_x$) increases with $n_e^2$. However, we note that a factor of 15 change in EM will correspond to less than a factor of 2 change in $T$ according to formula (2).

For the gas at temperature around $2-3 \times 10^6$ K the filling factor is, most of the time, seen to be very close to 1 (Dere 1982). However, small filling factors ($10^{-4} < \phi < 10^{-3}$) have been found for very hot plasma (up to $10^7$ K in Martens, Van Den Oord, & Hoyng 1985). The presence of small amounts of very hot plasma in nonflaring active regions of the corona ($T \geq 10^7$ K) has also been reported by Schadee, De Jager, & Svestka (1983), but no filling factor has been derived in this case. Unfortunately, these measurements had a low spatial resolution (pixel size is $8'' \times 8''$), so that it is difficult to get information about the spatial repartition of the hot gas in the loop. We will demonstrate below that these small filling factors for the high-temperature plasma represent a key feature, which indicates a strongly localized heating process.

2.2.4. Magnetic Field Strength and Electric Currents

Potential field extrapolations from photospheric magnetograms, which are a reasonable baseline assumption as mentioned earlier, provide a determination of the longitudinal component of the magnetic field (i.e., along the loop). The magnetic field strength is calculated to lie in the range $30 < B_{\text{GJ}} < 200$, where $l$ denotes the longitudinal direction. In order to extract some “typical coronal values” for $B$ we have used the values reported by Golub et al. (1980), supplemented with a magnetogram made the day of the 1989 September NIXT rocket flight.

Departures from the potential configuration can be deduced from vector magnetograms by comparing the extrapolated components with the observed ones (see Hargreaves et al. 1984 for more details on the technique employed). These authors observe evidence of large departures from the photospheric potential configuration: the largest of such “sheared magnetic fields," in terms of shear angle and spatial scale, are strongly related to the location of flare onset. For such nonpotential configurations, another component of the magnetic field exists in the azimuthal direction (the poloidal component $B_p$; see Fig. 3), which is characteristic of the presence of some currents flowing along the loop. We then have a nearly force-free configuration (i.e., $\mathbf{V} \times \mathbf{B} = \alpha \mathbf{B}$). Haisch et al. (1986) and Gary et al. (1987) have derived the quantity $\alpha$ throughout an AR. Results of such calculations show that, although $\alpha$ varies over the AR, a net current arches over the neutral line by flowing antiparallel to the magnetic field lines. Gary et al. (1987) give an average current density in the ARs of $4 \times 10^{-8}$ A cm$^{-2}$, while Haisch et al. (1986) give local peaks at $5 \times 10^{-7}$ to $10^{-6}$ A cm$^{-2}$, at a spatial resolution of about 2.5$''$. These currents flow from the negative polarity domain to the positive one.

![Fig. 3.—Cylindrical model of a current-carrying plasma column with a longitudinal magnetic field component $B_l$ and a poloidal component $B_p$.](image)
This is a very strong result which clearly indicates the presence of net currents in coronal loops. This is in opposition to the usual assumption that the total net current should be very low in coronal loops (see Parker 1990 and van Ballegooijen 1990, for example) which leads to models incorporating equally distributed currents along both directions.

2.2.5. Temporal Behavior

Coronal loops offer a variety of temporal behaviors (Sheeley & Golub 1979; Levine & Withbroe 1977). However, typical time scales for variability in individual AR loops are ~1 hr and changes in overall patterns occur over several days. Sheeley & Golub (1979) indicate that in the XBP, the individual loop structures vary with a time scale of ~6 minutes, which is about the cooling time of the plasma for their conditions.

Unfortunately, it is not possible from the observations to make direct inferences about the physical processes occurring in the corona: ionization and recombination time scales of tens to hundreds of seconds are involved in the appearance and disappearance of X-ray emission even if the temperature changes instantaneously (Golub, Hartquist, & Quillen 1989). The available data are consistent with the view that coronal loops brighten and fade on time scales comparable to those needed for formation of a given ionization state or for cooling of the plasma by radiation and conduction.

2.2.6. A Set of Typical Coronal Loops

The values which we adopt are summarized in Table 1 for various features: X-ray–bright points (XBP), active regions (AR), and large-scale structures (LSS). In the rest of the paper any result derived for typical coronal conditions for nonflaring ARs refers to the parameters in Table 2 describing three types of loops: a short dense one (type 1), a medium one (type 2), and a longer weaker one (type 3). For simplicity, the distribution of temperature along the loop and related flows are ignored. The usual assumptions of quasi-isotropy and isodensity for electrons and ions ($n_e \approx n_i$ and $T_e \approx T_i$) are made unless otherwise stated, so that the pressure is $p = n_e k T_e + n_i k T_i$, with $k$ the Boltzmann constant.

2.3. Rates of Energy Loss

Based on the loop parameters, we can derive the minimum heating rate necessary to sustain the coronal loops at such high temperature and density by considering what the losses are in this medium. There are mainly two classes of energy loss: the heating rate necessary to sustain the coronal loops at such high temperature and density by considering what the losses are in the Boltzmann constant.

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To obtain the optically thin part of the atmosphere, such as the transition region and the corona, the radiative losses take the form

$$L_r = n_e n_i P(T),$$

where the electron density $n_e$ and the ion density $n_i$ are usually assumed equal. $P(T)$ has been evaluated by a number of authors, and some fits based on the analytical approximation $P(T) = X T^4$ have been performed (see Rosner et al. 1978b or Priest 1982 for a more detailed discussion). The following form for $10^5 < T_K < 10^7$ will suffice for our purpose (Priest 1982):

$$P(T) = 10^{-18.666}T^{-1/2}.$$  

The radiative losses per unit volume then are

$$L_r = 10^{-18.666}T^{-1/2}n_e^2.$$  

$L_r$ is measured in [ergs cm$^{-3}$ s$^{-1}$]. Calculations of the conductive losses assume a spatially constant heat flux along the loop (Krall 1977). From this, the following formula is easily derived:

$$L_c = \frac{(5.14 \times 10^{-7}T^{7/2})}{(2L^2)}.$$  

Under the stationary conditions, the energy equation is simply a balance between the loss rate and the heating rate. The sum of the two loss terms thus gives the minimum heating rate necessary to sustain a loop in the solar corona. Numerical estimates for the set of typical coronal loops described in § 2.6 are displayed in Table 2.

Note that if an inhomogeneous loop is considered with, for instance hotter gas embedded in the loop, the calculation of the losses should include it through the filling factors of each population. Suppose, for example, that a loop has the bulk of its plasma at temperature $T_i$ and a small population at temperature $T_2 > T_i$ (the filling factor of $T_1$ being $\phi$), then the total losses for the loop will be

$$E_r = (1 - \phi)E_r(T_i) + \phi E_r(T_2)$$  

and

$$E_c = (1 - \phi)E_c(T_i) + \phi E_c(T_2).$$  

Since the filling factor of the hot gas is very small ($\phi \approx 10^{-4}$ to $10^{-3}$), this effect is negligible for the energy balance.

<table>
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<tr>
<th>Parameter</th>
<th>XBP (X-Ray Bright Point)</th>
<th>AR (Active Region)</th>
<th>LSS (Large-Scale Structure)</th>
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<td>0.04-0.01</td>
<td>0.01</td>
</tr>
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<td>2-3</td>
<td>2</td>
</tr>
<tr>
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<td>2-8</td>
<td>0.5</td>
</tr>
<tr>
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<td>$\phi E_r(T_2)$</td>
<td>$(1 - \phi)E_c(T_i) + \phi E_c(T_2)$</td>
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<tr>
<td>$E_r(T_2)$</td>
<td>$(1 - \phi)E_r(T_i) + \phi E_r(T_2)$</td>
<td>$E_c(T_i)$</td>
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3. HEATING THEORY

3.1. Basic Considerations

We propose a set of processes that are relatively simple and which should compare reasonably well with the available observations on the Sun. An effort is made to base our model on the detailed experimental and theoretical studies which have been carried on in terrestrial hot magnetized plasma work. We therefore can have some confidence in the reality of the processes which are discussed, and in the possibility that they may be occurring on the Sun.

From the observations summarized above, there is strong evidence that electrical currents are generated along the magnetic field lines of a coronal loop. These are studied in § 3.2. The precise mechanism by which currents are induced in the corona is not directly relevant to our model; what matters is only that field-aligned currents are present. The associated nonpotentiality of the magnetic field has indeed been measured by means of vector-magnetograms and leads one to reconstruct a current distribution in the loop consistent with our assumptions. The current densities thus generated experience relatively little magnetic diffusion and are confined at the periphery of the loop in a narrow, intermittent current sheath by the well-known skin effect. Here, "intermittent" refers to the phase preceding the recurring destruction of current sheaths by resistive instabilities. However, as we shall show, the stability of such a current sheath is a problem which also needs to be addressed.

Observations of coronal loops lasting for several hours suggest that they are macroscopically stable, which sets an upper bound on the total plasma current circulating in the loop. The high heating rate necessary to balance the radiative and conductive losses of the loop suggests that the electrical resistivity might be enhanced by several orders of magnitude relative to its classical value as studied by Tucker (1973) and Rosner et al. (1978a). We study this possibility in § 3.3. If the current sheath described above is so thin as to allow the current density to reach very high values and exceed a critical threshold, microinstabilities can be excited and lead to an anomalous electrical resistivity. This possibility is considered as a transitory stage within our model, during which the current density profile is peaked around the plasma axis is obtained. Although this is possible for the small typical dimensions of laboratory experiments, it is unrealistic with solar scales, as shown earlier. Figure 5 illustrates the skin current configuration in laboratory experiments (a) and as assumed in our model for coronal loops (b). Figure 6 describes the envisioned magnetic structure of the loop: the magnetic field lines are twisted in the outer current sheath, but stay essentially longitudinal in the loop's core where no current has penetrated. Spicer (1981) reaches the same conclusion in the context of solar flare models. We note that coronal loops in their entirety are seen as separate features, in a medium where the magnetic field is probably more or less homogeneous: the loop boundary originates at the photospheric level where the magnetic field is seen to be made of separate flux tubes (Title et al. 1989).

As discussed in § 2, coronal loops are seen to remain bright for several hours, often in a state that appears stationary. This macroscopic stability of the loops suggests that a constraint exists on the current density flowing along the loop. According
to a well-known criterion, a cylinder with an axial magnetic field \( B_r \), an azimuthal magnetic field component \( B_{\varphi} \), a length \( l \), and a radius \( a \) (see Fig. 3), becomes macroscopically unstable to a kink instability as soon as

\[
2\pi \left( \frac{B_l}{B_{\varphi}} \right) \left( \frac{a}{l} \right) \leq 1. \tag{11}
\]

This criterion has been well verified in laboratory experiments that include toroidal configurations. For a coronal loop with \( l = 2L \), it becomes \( B_{\varphi} > \pi(a/L)B_l \).

However, this criterion considers the ideal case of a free cylinder. As suggested by Hood & Priest (1979), a coronal loop is not free in the sense that the footpoints of the magnetic field lines are anchored in the photospheric layer, providing an effective stabilizing mechanism against macroscopic instabilities. They include this effect in their calculations and show that the critical pitch angle, \( \theta_{\text{crit}} = (B_{\varphi}/B_l)(2L/a) \), is a rapidly increasing function of \( L/a \) when \( L/a > 10 \), reaching the value \( 6\pi \). In our case, \( L/a = 25 \) so that we can take \( \theta_{\text{crit}} = 6\pi \), relaxing the criterion for the loop stability by a factor of \( 3 \)

\[
B_{\varphi} \leq 3\pi \left( \frac{a}{L} \right) B_l. \tag{12}
\]

For example, with \( a/L = 0.04 \) the loop is stable if \( B_{\varphi} < 0.37B_l \).

This constraint on the poloidal component of the field can in turn be expressed in the form of a constraint on the total plasma current flowing along the loop in a force-free configuration. The relation between the total current \( I \) and the poloidal magnetic field \( B_{\varphi} \) comes from the Maxwell equation \( \mathbf{\nabla} \times \mathbf{B} = 4\pi\mathbf{J} \), which leads to

\[
I = 5aB_{\varphi}, \tag{13}
\]

where \( a \) is in cm and \( B_{\varphi} \) in G. The constraint on the current then becomes:

\[
I \leq 15\pi a^2 \frac{B_l}{L}. \tag{14}
\]
We can evaluate the relevance of this new criterion by comparing the results of the current density averaged over the typical loops \( J \sim 1/(\pi a^2) \) given in Table 2 with the observations (whose spatial resolution is roughly the size of a loop): Haisch et al. (1986) give current density \( J_{\text{obs}} \sim 10^{-7} \) to \( 10^{-6} \) A cm\(^{-2}\) with a spatial resolution of 2.5. It appears that the calculated values are in the observed range, so that in the rest of this paper, we will assume this new criterion to be valid in coronal loops. Note that loops with either a stronger toroidal magnetic field or a smaller aspect ratio (compact loops) can stand higher currents, and then according to our model should appear brighter. We can deduce the local current density in the current sheath:

\[
J_c = \frac{I}{(\pi a^2 f)} ,
\]

(15)

where \( f \) is the filling factor of the current density in the loop. Substitution (14) in (15) yields

\[
J_c \leq 15 \frac{B_1}{fL} .
\]

(16)

### 3.3. Excitation of Microinstabilities and Anomalous Resistivity

Depending on plasma conditions, various microinstabilities, driven by the electron current, may take place; examples of these are the so-called Buneman, ion-cyclotron, ion-acoustic, and lower hybrid mode. Each of them is characterized by a threshold drift velocity \( v_d \) of the current carrying electrons. The relevant drift velocity threshold for each of them is, for instance, \( v_d > v_s \) for the Buneman instability, \( v_d > 0.3 v_s \) for the ion-cyclotron instability, and \( v_d > v_s \) for the ion-acoustic instability in a plasma with \( T_e > T_i \), where \( v_s \) is the thermal velocity of the electrons, and \( v_s = (kT/M)^{1/2} \) is the ion-sound velocity. The relevant numerical values are

\[
v_d = \frac{J_e}{n_e e} \approx 0.62 \times 10^{19} \frac{J_e}{n_e} \text{ cm s}^{-1} ,
\]

(17)

\[
v_e = (2kT_e/m)^{1/2} \sim 5.5 \times 10^5 T_e^{1/2} \text{ cm s}^{-1} ,
\]

(18)

\[
v_s = (kT_e/M)^{1/2} \sim 0.016 n_i^{1/2} \sim 10^4 T_i^{1/2} \text{ cm s}^{-1} ,
\]

(19)

where \( J_e \) in A cm\(^{-2}\) and \( n_e \) in cm\(^{-3}\) are the current and electron density in the current sheath electron, \( e \) is the charge in C and \( T_e \) is the electron temperature in K.

Rosner et al. (1978) have discussed in detail the possibility for the ion-sound instability to take place under conditions typical of coronal loops. The idea is that the electrons are preferentially heated (see Kaplan & Tsytovich 1973) and the \( T_e > T_i \) condition required for ion-acoustic turbulence is reached. The type of anomalous resistivity that has been proposed is of the form

\[
\eta_a = 4 \pi \frac{v_{\text{eff}}}{\omega_{pe}} (s)
\]

(20)

where

\[
\omega_{pe} = (4\pi n_e^2 m_e^{1/2})^{1/2} \approx 5.64 \times 10^4 n_i^{1/2} \text{ rad s}^{-1}
\]

(21)

and the effective collision frequency

\[
v_{\text{eff}} \approx \left( \frac{\omega_{pe}/100 v_d/v_e}{} \right) (s^{-1})
\]

(22)

(see Kaplan & Tsytovich 1973). Consequently

\[
\eta_a \approx \left( \frac{2/12\pi}{25} \right) \left( \frac{v_d}{v_e} \right) \omega_{pe}^{-1} (s).
\]

(23)

We note that for \( n \sim 3 \times 10^9 \) cm\(^{-3}\) and \( v_d/v_e \approx 0.016 \), we obtain \( v_{\text{eff}} = 7.1 \times 10^{15} \) s\(^{-1}\). The corresponding magnetic diffusion coefficient \( D_\phi = v_{\text{eff}}(e/\omega_{pe})^2 \) is about \( D_\phi = 7 \times 10^7 \) cm\(^2\) s\(^{-1}\). Thus the relevant anomaly factor \( D_\phi/D_\text{eff} \) is about \( 3.2 \times 10^4 \).

Given that the relevant threshold is \( v_d \geq v_s \), then

\[
J_e \geq n_e v_s
\]

(24)

where \( e \) is in C. We also have \( J_e = I/(2\pi n_a a) \), where \( a \) is the width of the current sheath in cm. We therefore find a constraint on the current sheath width

\[
\Delta a \leq 1/(n_e v_s 2\pi a) \text{ (cm)}.
\]

(25)

Table 3 shows calculated values for the ion-acoustic resistivity \( \eta_a \), the current sheath width \( \Delta a \), the current density in the sheath \( J_e \), and for the filling factor \( f = 2\Delta a/a \), relevant to the excitation of current-driven ion acoustic modes [eqs. (23), (24), (25)]. We point out that the filling factors of current density thus derived are substantially smaller than those observed by Martens et al. (1985) for the hot plasma.

The ohmic heating rate is given by

\[
E_{\text{oh}} = IU \text{ (ergs s}^{-1} \text{)} ,
\]

(26)

where \( U \) is the loop voltage and can also be expressed as

\[
U = 2L \eta J_e \text{ (statvolt)}.
\]

(27)

Since \( J_e = I/(2\pi n a a) \)

\[
E_{\text{oh}} = \frac{L \eta I^2}{\pi a a}.
\]

(28)

However we argue that the current density will rapidly diffuse over a layer \( \Delta a \) that is larger than that required by the condition \( v_d > v_s \). Given the smallness of the current layer needed for the onset of the relevant anomalous resistivity, we consider the onset of current-driven ion acoustic modes to occur at best only under transient conditions in order to allow the current density to distribute itself over a larger layer.

#### 3.4. The Dissipation of Magnetic Energy

3.4.1. A Reconnecting Instability-based Process

Relatively fast macroscopic instabilities driven by the current density gradients under a variety of conditions that can be close to or relatively far from ideal MHD marginal stability can be excited and involve magnetic reconnection processes associated with the finite electrical resistivity of the plasma (Coppi 1963; Coppi et al. 1976b) or the effects of finite electron inertia (Coppi 1964). When the current sheath is so thin that microinstabilities of the type studied in the previous section can develop, the ensuing resistive instability can broaden the skin current configuration over a width larger than the critical
threshold for the onset of ion sound instabilities. The plasma will return to a state where the d.c. electrical resistivity will be much lower. As we will show below, the destruction of the current layer then converts the energy stored in the poloidal magnetic configuration into particle heating and acceleration. We propose the following sequence of events for the loop heating cycle:

1. The plasma loop initially has a high electrical conductivity, and the electric field generated through the photospheric motions induces in turn a current in the loop outer layer over a distance for which the current density does not exceed the threshold of excitation of ion sound modes. (This may come possibly after a transitory phase during which these modes were excited, as described in § 3.3.)

2. The current sheath thus obtained can then be destroyed by current density gradient-driven instabilities, and magnetic energy converted into particle acceleration and plasma heating. During the lifetime of the skin-current configuration, a small filling factor of hotter gas is present.

3. Thermal energy transport across the magnetic field lines is locally enhanced by the excited instability, and heat diffuses across the loop. The input of thermal energy into the loop may lead to chromospheric evaporation if the time scales are adequate. As a consequence of high thermal conduction along magnetic field lines, the chromospheric material at the loop footpoints is heated up and evaporates into the corona (Kuin & Martens 1982). This leads to a density increase which induces brightness variability since the latter is proportional to $n_e^2$.

4. The resulting turbulence subsides and the cycle restarts with phase 1.

This sequence, involving heating through short-lived bursts, occurs repeatedly in a continual but not continuous process as long as the mechanism which generates the currents continues to operate. Therefore it is probable that both the filling factor and the temperature distribution throughout the loop volume vary with time following the sequence of events, which produces an intermittent heating process (see Fig. 7). Figure 8 illustrates a loop's brightness variability according to the described sequence. We note that there is experimental evidence of the fact that a well-confined plasma column tends to develop large-scale, unstable modes when the current is induced at a relatively fast rate and a skin current configuration begins to form; these are the so-called Mirnov oscillations (see Wesson 1986, for example).

The self-disruption of the current sheath involves a fast rate of thermal energy conversion and transport across the magnetic field lines, that redistributes the energy to a larger volume. In § 3.4.2 we discuss the characteristics of the disrupting modes. On the basis of observational results and energetic requirements of the loop, examples based on this model will be given in § 3.4.3.

### 3.4.2. Characteristics of Reconnecting Modes

We note that the dimensionless parameter that is used in dealing with the effects of the classical resistivity on modes producing magnetic reconnection is

$$\epsilon_n = \frac{D_m}{v_{Ap}\Delta a},$$

where $v_{Ap} = B_p/(4\pi n_v M)^{1/2} \approx 2 \times 10^9(3 \times 10^9/n)^{1/2}(B_p/50)$ (cm s$^{-1}$), so that

$$\epsilon_n \approx \left(\frac{50}{B_p}\right)\left(\frac{n}{3 \times 10^9}\right)^{1/2}\left(\frac{10^5}{\Delta a}\right)^{1/2}\left(\frac{3 \times 10^6}{T}\right) \times 10^{-10}.$$  (30)

This is so small that if we estimate the width $\delta_R$ of the region where magnetic reconnection is allowed to take place, as a resistive mode driven by the current density gradient is excited, this is of the order of

$$\delta_R \approx (\Delta a/10^5)(\epsilon_n/10^{-10})^{2/5} \times 10 \text{ (cm)}.$$  (31)

The growing rate of the resistive mode corresponding to the width of the reconnection layer (31) is of the order

$$\gamma \approx (\gamma_{Ap} \epsilon_n^{3/5})$$  (32)

with

$$\gamma_{Ap} \approx \frac{v_{Ap}}{\Delta a} \approx \left(\frac{B_p}{50}\right)\left(3 \times 10^9\right)^{1/2}\left(\frac{10^5}{\Delta a}\right) \times 5 \times 10^3 \text{ (s}^{-1})\text{).}$$  (33)

This instability produces magnetic islands (see Coppi & Friedland 1971), and we note that the linear theory of this mode breaks down when the size $\delta_i$ of the magnetic islands that it produces

$$\delta_i \approx (\Delta a)(B/B_p)^{1/2}$$  (34)
becomes of the order of
\[ \frac{\tilde{B}_r}{B_p} \approx \epsilon_\delta^{4/5} = 10^{-8}(e_\delta/10^{-10})^{4/5}, \tag{35} \]
where \( \tilde{B}_r \) is the perturbed radial magnetic field produced by the instability. The estimated value for \( \tilde{B}_r \) (\( \approx 10^{-8}B_p \)) from equation (35) is unrealistically small. Therefore we argue (see Coppi & Detragiache 1990) that a strictly linearized theory of the mode producing magnetic reconnection by the effects of the classical resistivity is not realistic, since the relevant magnetic field perturbation is so small, and adopt a theoretical model whereby a string of magnetic islands with a size
\[ \delta_{10} > \Delta \epsilon_\delta^{2/5} \tag{36} \]
is formed and becomes the relevant elementary distance instead of the reconnection layer \( \delta_R \) previously mentioned. Figure 9 gives a schematic representation of a loop in which magnetic islands have formed. We assume that in the presence of these small islands the electrons lose momentum at a rate corresponding to magnetic field diffusion represented by (see Coppi & Detragiache 1990)
\[ D_{10} \approx \alpha_D \left( \frac{v_e}{\Delta a} \right) \delta_{10}^2, \tag{37a} \]
where \( v_e \approx 10^9 \times (T_e/3 \times 10^6)^{1/2} \) (cm s\(^{-1}\)). Correspondingly,
\[ \epsilon_{\delta_{10}} = D_{10}v_e\Delta a \approx \alpha_D(\delta_{10}/\Delta a)^2(v_e/v_{Ap}). \tag{37b} \]
This state is unstable, with magnetic reconnection occurring over a distance of the order of (see appendix in Coppi & Friedland 1971)
\[ \delta_R \approx (\Delta a)\epsilon_\delta^{2/5} \approx \alpha_D \Delta a^{1/5} \delta_{10}^{4/5} (v_e/v_{Ap})^{2/5} \tag{38} \]
and a growth rate
\[ \gamma \approx (\delta_{10}/\Delta a)^{6/5}(\alpha_D v_e^{3/5}v_{Ap}^{2/5})/\Delta a \tag{39} \]
whose effect is to release the energy available in the current layer. Then we may argue that the instability will develop at this rate until the size of the magnetic islands it produces reaches a value of the order of
\[ \delta_{11} \approx \Delta a^{1/5} \delta_{10}^{4/5} (\alpha_D v_e/v_{Ap})^{2/5}. \tag{40} \]
In order to follow the mode evolution analytically, we may assume that when the level (40) is reached by the mode’s amplitude, the instability saturates and that a new equilibrium is reached. If the magnetic diffusion coefficient corresponding to the larger magnetic islands is
\[ D_{11} \approx \alpha_D(v_e/\Delta a)\delta_{11}, \tag{41} \]
a slightly faster instability develops until a new saturation stage is reached, etc.

We point out that in discussing the nature of the instabilities that can be excited, it is important to consider the propagation time of a magnetic field perturbation along the loop. This time

\[ \text{FIG. 8.—Qualitative illustration of the events causing loop brightness variability: (a) impulsive heating in a current sheath at the loop periphery; } T \text{ is locally } \approx 10^7 \text{ K;} (b) \text{ heat diffuses throughout the loop, raising its temperature to } \approx 10^6 \text{ K;} (c) \text{ heat conduction downward leads to chromospheric evaporation and thereby to a density rise in the loop and a brightness enhancement; and (d) when heating rate falls below energy loss rate material condenses back to the chromosphere; } T \text{ and } n \text{ decrease and cycle restarts at (a).} \]
is not usually relevant in laboratory plasmas where the Alfvén propagation time is shorter than all other characteristic times. In the present case we shall assume that the instability that destroys the current layer does not develop coherently all along the loop as a normal mode, but that it will affect regions of limited length. The maximum value $L_y$ can be estimated as

$$L_y \approx y v_P,$$  \hspace{1cm} (42)

where $y$ is the instability growth rate. It is clear from these considerations that the instability is then to be treated as a nonnormal mode (Coppi et al. 1966), localized over a distance $L_y$ along the magnetic field lines.

Another aspect to be considered is that simultaneous strings of magnetic islands can be assumed to be formed on a sequence of magnetic surfaces at different radii of the plasma column, within the distance $\Delta a$. Therefore magnetic reconnection can occur at the same time over a substantial portion of the current density profile.

3.4.3. Comparison with Observations and Energy Balance

The observed filling factors of hot gas ($\phi = 10^{-4} - 10^{-3}$, Martens et al. 1985) give an indication of the current sheath size. If the filling factor for the hot gas is of the order of the current density, $\phi = f$, then we have $\phi = 2\Delta a/a$. Once $\Delta a$ is known, an upper bound for the time scale of the disruption can be derived, using equation (39) in the limit corresponding to the slowest rate. Examples of current sheath widths and relevant time scales upper bounds for each of the typical loops are summarized in Table 4. The process can be viewed as evolving through short-lived bursts. This is in agreement with the observations of Sheeley & Golub (1979) who find that the typical time scale for the brightness variability is of the order of the loop cooling time and with the Parker (1988) suggestion of quasi-impulsive events.

The energy density that can be converted into flow and thermal energy, as a result of the current layer instability involving magnetic reconnection, can be assumed to be of the order of

$$fB^2_{\phi}/(8\pi)$$  \hspace{1cm} (43)

(see Coppi et al. 1976a), where $f$ is the current filling factor $f = 2\Delta a/a$. The transformation of magnetic energy into thermal energy was discussed in Coppi & Friedland (1971). They give a quadratic form in the amplitude of a mode producing magnetic reconnection (see eq. [A15] of their paper) whose components can be related to the various forms of energy variation involved. They show that “the energy extracted from the magnetic field goes into kinetic and ohmic dissipation energy.” Thus, the heating rate for each event is the ratio between the energy converted by one event and the time-scale of this event

$$L_h \approx fB^2_{\phi}/(8\pi \tau_c)$$  \hspace{1cm} (44)

where $\tau_c$ is the typical disruption time of the current sheath.

Table 4 presents minimum values for the heating rate according to equation (44), when the plasma filling factor $\phi$ is assumed to be around $10^{-3}$ and where $\tau_c$ is the upper bound previously derived. Also the duty cycle, which is the ratio between the energetic requirements of the loop and the heating rate provided by our model (the percentage of the time taken by the energy release process) is given. The rest of the time is used for the transverse thermal transport and the current sheath regeneration. We point out that, given the relatively short current sheath lifetime, the total plasma current in the loop in practice cannot rise to the point of reaching the threshold for the onset of an ideal MHD kink mode.

These large values for the heating rate (and consequently small duty cycles) suggest that coronal heating in ARs is not a simple continuous mechanism. In fact, we should speak of “a number of dynamical processes” intermingled. Disruption of the current sheath and redistribution of the heat occur on the same time-scale (10–100 s) as the conduction of energy along the magnetic field lines leading to chromospheric evaporation and the atomic processes which lead to the formation of observable line emission. Thus a detailed calculation of the observable signatures of the proposed mechanism would be very complex, as illustrated in Figure 10.
In the proposed configuration, the question of the radial transport (i.e., across the magnetic field lines) still needs to be addressed. A detailed treatment of this issue is beyond the scope of this paper, and we may limit ourselves to a few considerations.

The self-disruption of the magnetic configuration provides a large enhancement of the particle crossfield conduction given that the magnetic field lines undergo reconnection. Foukal (1978) points out that observations of cool plasma flows in coronal loops are difficult to understand “unless reconnection allows plasma motion across field lines.”

A continuous state of magnetic turbulence may exist inside coronal loops, possibly in relation with the footpoint motions. This could increase the transverse heat diffusion. In fact, an enhanced cross-field thermal conductivity, relative to that predicted by the collisional theory, is commonly observed in laboratory experiments and appears to be associated with the excitation of various kinds of fluctuations.

Although the proposed sequence only mentions a single current layer, the finite lifetime of current sheaths, together with the requirement that the coronal portion of a loop must remain in torsional equilibrium with its photospheric footpoints (see Rosner et al. 1978a), makes likely the presence of a number of current sheaths at various stages of evolution, radially dispersed throughout the loop, continually forming and disrupting. This would greatly relax the requirements for heat transport across the loop.

4. SUMMARY AND CONCLUSIONS

We have considered recent observational advances such as coronal loop geometry deduced from high spatial resolution X-ray images (Golub et al. 1990), the presence of high electric currents flowing along coronal loops as determined from potential extrapolations of vector-magnetograms (Haisch et al. 1986), and measurements of small filling factors ($\sim 10^{-3}$) of very hot gas ($\sim 10^7$ K) in coronal loops (Martens et al. 1985) which indicate the existence of localized heating processes. The new observations have been combined with well-established theoretical results from laboratory experiments in order to discuss possible models for heating of coronal loops based on electric current dissipation.

First, we summarize the observational basis of our study, from which we deduce a set of three typical coronal loops that are used in the numerical estimates. We argue, without specifying a mechanism but by relying on the observational data, that field-aligned currents are generated along the coronal loops. Because of the high electrical conductivity in the loop, a skin-effect confines the induced currents into thin annulus-shaped “current sheaths.” The macroscopic, ideal MHD, stability of the loop requires that the total electric current is limited by an upper bound which is found to be in the same range as the observed currents.

We first note that the current sheath width could be sufficiently small during a relatively fast transient to allow the current density to reach a value high enough to trigger microinstabilities such as ion-acoustic turbulence, which in turn generates an anomalous electrical resistivity (cf. Rosner et al. 1978a). For this case we argue that the width of the current-carrying layer would be broadened rapidly without being disrupted.

Then we propose that the excitation of a disruptive instability, driven by the current density gradient and assumed to occur for values of the current density well below the threshold for the onset of known microinstabilities such as ion-acoustic modes, converts the energy associated with the current-related magnetic field (the poloidal component $B_p$) into particle acceleration and thermal energy. The resulting heating rate is estimated to exceed considerably the energetic requirements of the loop. Thus, the overall heating process can be envisioned to proceed by short bursts, corresponding to intermittent disruptions and regenerations of a current sheath configuration.

A strong intermittent heating rate may involve an evaporation mechanism, followed by a condensation mechanism, which results in density variations. This is consistent with the observed rapid brightness variability of coronal loops (see Sheeley & Golub 1979). It is also consistent with the observation of low-level impulsive hard X-ray bursts in active regions under conditions which would not normally be called flare-like (Lin et al. 1984).

Because of the relatively low transverse thermal conduction, initially only a small fraction of the loop volume is heated and exhibits a much higher temperature than the average value. This is in agreement with the observations of small filling factors of very hot gas in coronal loops (Martens et al. 1985). When the self-disruption of the current sheath occurs, the thermal energy is rapidly dispersed into a larger volume. Throughout this paper, numerical estimates for a set of three typical coronal loops are used to illustrate the proposed model.

On the basis of this model we envisage the following sequence:

1. The plasma loop initially has a high electrical conductivity, and the electric field generated through the photospheric motions induces in turn a skin-current in an outer layer of the loop.

2. The current sheath thus obtained is destroyed and magnetic energy is converted into particle acceleration and plasma heating. During the lifetime of the skin current configuration, a small filling factor of hotter gas may be observed, since the transverse thermal conductivity is relatively low.
3. Thermal energy transport across the magnetic field lines is locally enhanced by disruption of the current sheath, and heat diffuses across the loop. This input of energy may lead to an evaporation process if the time scales are adequate: as a consequence of high thermal conduction along magnetic field lines the chromospheric material at the loop footpoints is heated up, and evaporates into the corona (Kuin & Martens 1982). This leads to a density increase which induces brightness variability since the latter is proportional to n². The resulting turbulence subsides and the cycle restarts with phase 1.

Dynamical processes which are all simultaneously happening on similar time scales are conduction of energy along the magnetic field lines leading to chromospheric evaporation; disruption of the current sheath and redistribution of the thermal energy; and formation of observable line emission by microscopic atomic processes. Because all these processes occur on time scales of 10–100 s, a detailed evaluation of the observable signature of the proposed sequence is difficult to carry out. High-resolution X-ray observations of the corona, photospheric vector-magnetograms, and emission measure analysis, combined with spectroscopic measurements in order to detect small amounts of very hot plasma in coronal loops, and the relevant filling factors, will be necessary in order to carry the model we have proposed to a further degree of development.

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REFERENCES

Little, S. J., & Krieger, A. S. 1977, BAAS, 9, 341

Priest, E. R. 1982, Solar Magnetohydrodynamics (Dordrecht: Reidel)
Spitzer, L. 1962, Physics of Fully Ionized Gases (New York: Interscience)