NUMERICAL SIMULATIONS OF FLARES ON M DWARF STARS. I. HYDRODYNAMICS AND CORONAL X-RAY EMISSION

CHUNG-CHIEH CHENG
E. O. Hulburt Center for Space Research, Naval Research Laboratory, Washington, DC 20375

AND

ROBERTO PALLAVICINI
Osservatorio Astrofísico di Arcetri, Largo Enrico Fermi 5, Florence, Italy

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ABSTRACT

In order to understand the large variety of physical parameters observed in stellar X-ray flares, and their time evolution, we have made a series of numerical simulations of flare loop models with different values of loop size, flare energy input, and initial loop conditions. Our purpose is to study in detail the physics of flare hydrodynamic processes and to determine how the hydrodynamic evolution varies in parameter space. Our model solves the full set of mass, energy, and momentum equations in a magnetically confined loop structure which extends from the chromosphere to the corona. The computed results involve important physical processes such as the expansion of heated coronal gas, the evaporation of chromospheric material, the compression of the plasma at the loop footpoints, the gradual and/or catastrophic cooling of the gas during the decay phase. We also calculated the expected X-ray emission and compared it with X-ray light curves obtained by the EXOSAT satellite. Our simulations can reproduce the general characteristics of the observed X-ray flares, including their intensities, time profiles, and average coronal temperatures. We find that it is possible to produce X-ray emission measures of up to \( \sim 10^{52} \) cm\(^{-3} \) with fairly small loops with a total length of only \( 2 \times 10^8 \) cm. Small- and medium-sized flares can thus be produced in short loops, in contrast to previous conclusions that required large loops, of the size of a stellar radius, even for small flares. We find, however, that the amount of material that can be evaporated from the chromosphere does not vary linearly with the flare energy input, thus making increasingly more difficult to obtain large emission measures by simply increasing the energy input in a small loop. Under physically realistic conditions, large loops with length of \( 8 \times 10^9 \) cm are required to produce the largest emission measures (\( \sim 10^{-3} \) cm\(^{-3} \)) observed in stellar flares. The large evaporation velocities of more than \( 400 \) km s\(^{-1} \) predicted by our models should be observable as blue shifts in X-ray line profiles with the high resolution spectrometers planned for future X-ray missions such as AXAF.

Subject headings: hydrodynamics — stars: coronae — stars: flare — stars: late-type — stars: X-rays

1. INTRODUCTION

During the past decade, with the advent of space observations at X-ray and UV wavelengths, much progress has been made in the study of the outer atmospheres of cool stars. These observations have strongly suggested that the atmospheres of late-type stars are remarkably similar to that of the Sun, though often much more energetic. As is well known from high-resolution XUV and X-ray images obtained from Skylab, the outer atmosphere of the Sun is dominated by magnetic field structures, with closed loops in active regions and open field lines in coronal holes (Vaiana et al. 1973; Tousey et al. 1973; Underwood et al. 1976). Recent observations have shown that magnetic fields, starspots, and active regions also exist on many late-type stars (Bopp & Evans 1973; Vogt 1975; Eaton & Hall 1979; Marcy 1984; Poe & Eaton 1985; Giampapa 1985; Saar & Linsky 1985; Saar 1987, 1990). It is therefore plausible to suppose that the outer atmospheres of these stars are magnetically structured like the Sun, with closed loops being the dominant features in active regions.

The realization that the atmospheres of late-type stars are solar-like has stimulated efforts for interpreting their atmospheric structures in terms of magnetic loop models. These models extended from the solar case, give a fairly good account of the observed properties of cool stars (Linsky 1983; Antiochos & Noci 1986; Stern, Antiochos, & Harnden 1986). Flares observed from M dwarf stars have also been interpreted, in analogy to solar flares, as produced by the sudden heating of magnetic loops (cf. Haisch 1983). Most flare models developed so far, however, are only order-of-magnitude estimates of the relevant physical parameters, such as size, density, temperature, and cooling times. Although they are useful in establishing the first-order physics of a stellar flare, they cannot elucidate the details of the hydrodynamic evolution and energy transfer processes in the flaring plasma.

High-resolution XUV and X-ray observations (Widing & Cheng 1974; Cheng & Widing 1975; Pallavicini, Serio, & Vaiana 1977) have conclusively demonstrated that solar flares occur in loops that connect the low-temperature (\( \sim 10^4 \) K) chromospheric layers to the corona (where the temperature during the flare can reach values as high as \( 10^6 \) K). Simultaneous observations over a wide range of wavelengths, from hard X-rays to the radio, have further demonstrated that flares are very complex phenomena, involving many interrelated physical processes. In order to understand these complicated processes, numerical models have been developed by many authors and applied to the interpretation of the solar data (Nagai 1980; Cheng et al. 1983; Pallavicini et al. 1983; MacNeice et al. 1984; Cheng, Karpen, & Doschek 1985; Fisher, Canfield, & McClymont 1985; Peres et al. 1987). These numerical simulations have greatly helped in elucidating the
many hydrodynamic phenomena that result from the heating of a flaring loop such as, for instance, the propagation of high-speed conduction fronts and shocks, the formation of compressed plasma at the loop footpoints, the evaporation of the heated chromospheric material into the coronal part of the loop, and the cooling and condensation of plasma during the flare decay. As a result of these numerical simulations, observations of phenomena such as blueshifted components in high-temperature Ca xix and Fe xxv lines, downflows in lower temperature spectral lines, and enhanced X-ray and UV emission in solar flares can now be understood on a more firm physical basis.

Recent X-ray observations of stellar flares from M dwarf stars with good time resolution and adequate spectral coverage, particularly those obtained by the European satellite EXOSAT, have shown that stellar flares resemble closely solar flares, although on a much larger energy scale (see, e.g., Pallavicini, Tagliaferri, & Stella 1990). Observations of stellar flares still lack the fine details available for solar flares, yet the large body of radio, optical, UV, and X-ray observations now available (cf. Byrne 1989) justify a more detailed study of stellar flare models. Previously, Katsova, Kosovichev, & Livshits (1981), and Livshits et al. (1981) used a gas dynamic model to explain the optical emission of stellar flares as due to the compression of high density chromospheric plasma. Reale et al. (1988) simulated the hydrodynamic behavior of a specific X-ray flare observed by the Einstein satellite on the dMe star Proxima Centauri. They concluded that in order to account for the time evolution of the observed X-ray emission, the flare must have occurred in a loop with a length nearly as large as the stellar radius, in contrast to what is typically observed for the Sun. Fisher & Hawley (1990) and Kopp & Poletto (1990), on the other hand, used simplified analytical approximations to the hydrodynamic equations to predict the expected evolution of stellar X-ray flares. Their results, however, have to be checked against the more accurate numerical calculations in order to verify the plausibility of their assumptions.

In this paper, we extend our previous work on numerical simulations of solar flares to the stellar case (for previous solar work, see, e.g., Cheng et al. 1983 and Pallavicini et al. 1983). As our main purpose is to understand the basic physics of stellar flares, and the large variety of observed parameters, we adopt the approach of making a series of numerical simulations of flare loop models with different values of loop size, flare energy input, and initial loop conditions, rather than simulating a particularly well observed flare as was done by Reale et al. (1988). Specifically, our aims are (1) to see how the flare energetics and physical parameters vary with flare loop size, energy input, and initial loop conditions; (2) to determine the temporal evolution of the physical parameters of a flaring loop and to study the dynamic evolution of the coronal plasma during the flare; and (3) to calculate from the hydrodynamic results the expected X-ray spectroscopic signatures such as X-ray emission measures, X-ray line intensities, and line profiles. It is important to note that X-ray observatories of the next generation, such as the Advanced X-Ray Astrophysics Facility (AXAF), will have instruments capable of observing stellar X-ray line profiles with sufficiently high sensitivity and spectral resolution. It should be possible, therefore, in the near future, to compare the spectroscopic signatures predicted by our models with X-ray observations of stellar flares.

In the following sections, we first describe a few observations of stellar flares that are relevant for the choice of the input parameters for our simulations. In §3, we describe the hydrodynamic code and the numerical techniques used. In §4, we present the hydrodynamic results and, in §5, the predicted X-ray light curves and spectra. Finally, we discuss our results in §6.

2. OBSERVATIONAL CONSTRAINTS

Although flares have been detected from a large variety of stars with convective envelopes (see, e.g., Pettersen 1989), by far the most common class of flaring objects are the UV Ceti-type flare stars. These stars are typically of spectral type dMe. Flares from UV Ceti stars have been observed at radio, optical, UV, and X-ray wavelengths (cf. Haisch & Rodonô 1989 and references therein). The observations in X-rays, from which most of our knowledge of the high-temperature component of stellar flares is derived, have been reviewed comprehensively by Pallavicini et al. (1990) who used a large body of data from the European X-ray satellite EXOSAT. In this section we briefly describe the observed properties of dMe stars and their flares, with the purpose of selecting the input parameters for our numerical models.

The dMe stars, of which the UV Ceti-type stars are a subclass, are cool red dwarfs, with effective temperatures around 3000 K and hydrogen Balmer lines in emission. Many UV Ceti stars show quasi-periodic variations in the optical light curve, which are attributed to dark spots on their surface. Recent measurements show the existence of magnetic fields of several kilogauss covering up to 80% of the stellar surface (Saar 1987, 1990). Observations in the UV have shown that the outer atmosphere of these stars strongly resembles that on the Sun, with the existence of a hot chromosphere and transition region. Atmospheric models indicate that the temperature of the chromosphere at the base of the transition region ranges from 4300 to 25,000 K (Rodonô 1986) and that the chromosphere density is somewhat higher than for the Sun. A typical temperature of \(\sim 10^4\) K and a density of a factor of 10 larger than that for the Sun can be taken as representative for the chromospheres of dMe stars (Cram & Mullan 1979).

Quiescent X-ray emission from dMe stars varies from star to star but indicates in all cases the presence of a corona with a temperature in the range 3 \(\times\) 10^6 to \(\sim 1\) \(\times\) 10^7 K (Swank & Johnson 1982; Haisch et al. 1983). In analogy with the solar case, the high-temperature coronal plasma is believed to be confined in magnetic loop structures. The quiescent X-ray emission is rather steady with variations of less than a factor 2 over intervals of several months (Pallavicini et al. 1990). These fluctuations are most likely due to the emergence of new magnetic regions, or to decay of old active regions. X-ray flare activity appears to be rather common, with typically one flare for each \(\sim 10\) hr continuous EXOSAT observation. About 32 flares were observed by the Low Energy (LE) experiment on EXOSAT in the spectral band 0.05–2 keV, but only a few of them were observed by the less sensitive Medium Energy (ME)
experiment in the spectral band 1–10 keV. Figure 1 shows a few examples of flares observed by EXOSAT from UV Ceti-type stars. As can be seen from the figure, the time development of these flares is similar to that of solar flares, with a fast rise and a more gradual decay, and time scales ranging from several minutes to several tens of minutes.

Analysis of low-resolution ME spectra indicates flare temperatures in the range \( \sim 2 \times 10^7 \) to \( \sim 4 \times 10^7 \) K. The temperature appears to decrease during the flare decay, indicating cooling of the plasma. Note that these are average temperatures obtained by integrating the 1–10 keV ME flux over a cooling of the plasma. Note that these are average temperatures obtained by integrating the 1–10 keV ME flux over a substantial part of the flare evolution (sometimes even throughout the entire flare lifetime owing to the low observed ME count rate).

The peak temperatures reached during the flare may thus be somewhat higher than the quoted values. The flare emission measures [EM] derived from LE and ME data range from \( \sim 10^{51} \) to \( \sim 10^{53} \) cm\(^{-3}\). The total X-ray energies radiated throughout the flare lifetime range from \( \sim 3 \times 10^{39} \) to \( \sim 10^{42} \) ergs. Thus, while the observed time scales and average temperatures are similar to those of solar flares, the released energies and [EM]'s are, in general, much larger than those in solar flares. However, the strongest solar flare is as intense as the 1980 flare on the dMe star Proxima Centauri observed by the Einstein Observatory (Haisch et al. 1983; Haisch, Strong, & Rodonô 1991). In spite of the difference in energy scales, the EXOSAT observations clearly show that the phenomenology of stellar flares is similar to that of solar flares. In the remainder of this paper, we will assume therefore that the basic physical processes of energy release and transfer are also similar in solar and stellar flares.

There is evidence from the EXOSAT data, that two types of flares may exist on dMe stars: (1) impulsive flares, with rise times of a few minutes and decay times of tens of minutes; and (2) long duration flares, with lifetimes of the order of hours. The impulsive stellar flares are reminiscent of compact solar flares, which are associated with confined magnetic loops whose topology remains unchanged throughout the flare evolution. On the contrary, long-duration stellar flares are more similar to solar two-ribbon flares, which are believed to originate in magnetic structures that relax from an open configuration to a closed one (Pallavicini, Serio, & Vaiana 1977; Poletto, Pallavicini, & Kopp 1988). The loop models we simulate in this paper apply only to the more common class of impulsive stellar flares on dMe stars, and we do not consider the more energetic, longer duration flares typically observed from RS CVn binaries (Pallavicini & Tagliaferri 1989).

3. NUMERICAL MODEL

As mentioned above, our purpose in this paper is to investigate stellar flare simulations in the parameter space. We assume that stellar flares occur in a single magnetically confined loop structure as suggested by their analogy with solar flares. Although the atmospheric conditions on dMe stars can vary from star to star, they are close enough that we can adopt an average model to represent the typical conditions on a dMe star. For our calculations we have assumed a surface gravity twice that of the Sun. We have done also some simulations with a surface gravity 3 times that of the Sun. The results of using 3 times the solar gravity do not differ too much from the case of using twice the solar gravity. The length of the loop has been assumed to be either \( 2 \times 10^9 \) or \( 8 \times 10^9 \) cm. Both these values are much smaller than the radius of a typical dMe star, in analogy to what is observed for solar compact flares.

The geometry of our model is shown schematically in Figure 2. We consider a semicircular loop of constant cross section with the footpoints anchored in the dense chromospheric layers. The coordinate \( z \) is defined as the distance from the base of the loop measured along the loop itself. The temperature in the chromospheric portion of the loop is set equal to a constant value of 9000 K, as suggested by observations and models. The plasma in the loop is treated as a two-component fluid composed of electrons and ions. The plasma consists primarily of fully ionized hydrogen and helium, with a helium number density equal to 6.3% of the hydrogen number density (Ross & Aller 1976). We assume therefore that the elemental abundances in dMe stars are similar to those of the Sun. To obtain the evolution of the flare loop parameters, we solve the full set of time-dependent two-fluid hydrodynamic equations of mass, momentum, and energy conservation. The equations have been presented in Cheng et al. (1983) and will not be repeated here.

In the energy equation, we use the optically thin radiative loss function \( \phi(T_e) \) calculated as in Raymond & Smith (1977), where \( T_e \) is the electron temperature of the emitting plasma. For temperatures below \( 10^5 \) K, we use the semiempirical loss function \( \phi(T_e) = 8 \times 10^{-36} T_e^3 \), as given by McClymont & Canfield (1983). Below 9000 K, the radiative losses are assumed to be zero. The cutoff of optically thin radiative losses at this temperature is meant to mimic the trapping of radiation in the optically thick chromosphere layers. Since we are primarily

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interested in the general hydrodynamic response and in the high-temperature coronal plasma, the use of optically thin radiative losses is sufficient for our purpose. Although the chromosphere is not adequate to describe the physical conditions and radiative transfer in the chromosphere, it does not, however, affect significantly the time period of exponential decay with a time constant $r$. The spatial dependence of the heating function is assumed to be a Gaussian.

For the conductive term, the Spitzer conductivity has been used. It is well known that when the electron mean free path is greater than the temperature scale height, the classical Spitzer conduction breaks down, and the heat flux becomes saturated and delocalized. In these conditions, the heat flux depends on the global density and temperature structure of the plasma, in contrast to the classical case in which the heat flux is determined locally. The situation is especially severe in the initial phase of flare heating when the temperature gradient is extremely steep and the loop density is relatively low. Karpen & DeVore (1987) have studied the effects of nonclassical conductivity in hydrodynamic flare simulations, using an integral formulation that takes into account the nonlocal effects in the heat flux. Their results show that the nonlocal conductivity leads to preheating of chromospheric material and to evaporation which starts earlier and is less intense than in the classical case. The difference in temperature structure of the heated flare loop between the classical and nonclassical cases, however, becomes less and less noticeable as the flare heating progresses, owing to the increasing density in the loop.

Although the nonclassical heat conductivity affects importantly the energy transport in the flare loop, particularly in the initial phase, the hydrodynamic response of the loop is the same as for the classical conductivity. Also the total evaporated material into the coronal part of the loop is about the same in both cases. Since in this paper we are mainly concerned with the overall hydrodynamics of the flare and the resultant X-ray emission, the use of classical conductivity will not result in any loss of important physics and will not alter any of our conclusions. Hydrodynamic codes that include the nonlocal heat conductivity are available to us, but their use is prohibitively expensive and impractical for the purpose of simulating many stellar flares in parameter space.

The coupled hydrodynamic equations with proper boundary and initial conditions are solved numerically in an Eulerian scheme with an adaptive grid. The numerical algorithm used for solving the convective parts of the equations is the flux-corrected transport (FCT) method developed by Boris & Book (1976), while the nonconvective parts are solved implicitly. Time-splitting techniques are used for coupling time advances of terms representing various physical processes, such as radiation and conduction (Oran & Boris 1981). The FCT algorithm is especially suited to treat problems with steep gradients, such as shocks and the transition region, which arise in our flare calculations. Since we use a symmetrical loop and a flare heating function that is also spatially symmetrical and applied at the top of the loop, we need only to consider one-half of the loop (see Fig. 2). The length of the half-loop is denoted as $Z_{\text{max}}$ and includes the chromosphere, transition region, and coronal part. A total of 200 grid points are used to cover the half-loop and the grid size varies from the chromosphere to the corona, with the transition region covered by cells of the smallest size of about 10 km. Since a moving adaptive grid is used, the transition region is always covered by the smallest size cells during the flare evolution. In this scheme, the grid mesh is regridded at each time step so as to follow the motion of the transition region (cf. Oran & Boris 1981).

We have constructed three initial loop models with different loop lengths and equilibrium temperatures and densities as listed in Table 1. These loops are static equilibrium loops obtained by solving the hydrodynamic equations for given values of loop length and pressure at the base of the transition region. The boundary condition at the loop top is specified by setting the temperature gradient there to be zero. A constant quiescent loop heating rate ($S_0 \text{ ergs cm}^{-3} \text{ s}^{-1}$) is obtained such that it satisfies the specified base pressure and loop length. We have used two loop lengths, one with a semilength $Z_{\text{max}} = 1.45 \times 10^8 \text{ cm}$ and the other with $Z_{\text{max}} = 4.65 \times 10^9 \text{ cm}$, respectively. Since we assume for all models a chromosphere which extends up to $z = 5 \times 10^8 \text{ cm}$ at the loop footpoints, the total length for the coronal part of the loop is roughly $2 \times 10^9$ and $8 \times 10^9 \text{ cm}$, respectively, for the two cases.

Table 1 gives the initial loop top temperatures and densities. Models A and B are short loops with the same loop length but different base pressures that resulted in different loop top temperatures and densities; one with temperature of $2.7 \times 10^7 \text{ K}$ and density $1.5 \times 10^{10} \text{ cm}^{-3}$ (model A) and the other with temperature of $1.1 \times 10^7 \text{ K}$ and density of $1.3 \times 10^{11} \text{ cm}^{-3}$.

<table>
<thead>
<tr>
<th>Model</th>
<th>Semilength</th>
<th>Coronal Semilength</th>
<th>Base Pressure</th>
<th>Loop Top Temperature</th>
<th>Loop Top Density</th>
</tr>
</thead>
<tbody>
<tr>
<td>A.....</td>
<td>$1.45 \times 10^9$</td>
<td>$1 \times 10^9$</td>
<td>11</td>
<td>$2.7 \times 10^9$</td>
<td>$1.5 \times 10^{10}$</td>
</tr>
<tr>
<td>B.....</td>
<td>$1.45 \times 10^9$</td>
<td>$1 \times 10^9$</td>
<td>384</td>
<td>$1.1 \times 10^7$</td>
<td>$1.3 \times 10^{11}$</td>
</tr>
<tr>
<td>C.....</td>
<td>$4.65 \times 10^9$</td>
<td>$4 \times 10^9$</td>
<td>131</td>
<td>$1.1 \times 10^7$</td>
<td>$3.7 \times 10^{10}$</td>
</tr>
</tbody>
</table>
loops are all much smaller than the pressure scale height, their heating parameters. Since we know very little about the details of a static loop with higher loop temperature can be obtained only with a higher initial base pressure. Rosner, Tucker, & Vaiana (1978). Note that for a given length, and function are schematically shown in Figure 2. The flare heating during the flare also includes the quiescent heating, which is, however, many orders of magnitude less than the flare heating. The input parameters used for the 10 flare simulations are summarized in Table 2. The total flare energy deposited in the loop to the sudden deposition of energy is similar for all the model parameters. In this case the flare energy input drops to zero at the end of the time interval $t = 300$ s. Then we consider the case when there is additional heating during the decay phase, with the heating rate decreasing exponentially to zero. As we shall see later, these two cases have drastically different time evolutions thus providing an important diagnostic tool for the flare decay phase.

4. HYDRODYNAMIC RESULTS

We find that, although the basic hydrodynamic response of the loop to the sudden deposition of energy is similar for all the models we have run, there are, however, important quantitative differences in their evolution. We first present the general hydrodynamic results and then discuss the effect of changing the model parameters.

In order to illustrate the main hydrodynamic processes during the flare evolution, we show in Figures 3, 4, and 5 the spatial distributions of electron temperature, pressure, and velocity at different times for model 3 (short loop), model 8 (large loop), and model 9 (short loop). As can be seen from these figures, the hydrodynamic responses are rather similar and can be conveniently described in terms of the following evolutionary phases.

| TABLE 2 |
|---|---|---|---|---|---|---|
| Model | $L^e$ (cm) | $\sigma$ (cm) | $S_{\rho0}$ (ergs cm$^{-3}$ s$^{-1}$) | $t_1$ (s) | $t_2$ (s) | Preflare Model |
| (1) | (2) | (3) | (4) | (5) | (6) | (7) |
| 1 | 1 x 10$^9$ | 3.0 x 10$^8$ | 300 | 100 | 500 | A |
| 2 | 1 x 10$^9$ | 2.5 x 10$^8$ | 400 | 300 | 0 | A |
| 3 | 1 x 10$^9$ | 2.5 x 10$^8$ | 500 | 300 | 0,200,500 | A |
| 4 | 1 x 10$^9$ | 2.5 x 10$^8$ | 500 | 300 | 0,200 | B |
| 5 | 4 x 10$^9$ | 3.0 x 10$^8$ | 500 | 300 | 0,400 | C |
| 6 | 4 x 10$^9$ | 2.5 x 10$^8$ | 500 | 300 | 0 | C |
| 7 | 4 x 10$^9$ | 5.0 x 10$^8$ | 150 | 300 | 0,200 | C |
| 8 | 4 x 10$^9$ | 3.0 x 10$^8$ | 500 | 300 | 0,400 | C |
| 9 | 1 x 10$^9$ | 2.5 x 10$^8$ | 500 | 300 | 0,400 | A |
| 10 | 1 x 10$^9$ | 2.5 x 10$^8$ | 500 | 300 | 0,400 | A |

* Coronal semilength.

Generally, we run two cases for each model. First we assume that there is no heating in the flare decay (i.e., $t_2 = 0$). In this case the flare energy input drops to zero at the end of the time interval $t = 300$ s. Then we consider the case when there is additional heating during the decay phase, with the heating rate decreasing exponentially to zero. As we shall see later, these two cases have drastically different time evolutions thus providing an important diagnostic tool for the flare decay phase.

### TABLE 3

| DERIVED FLARE PARAMETERS |
|---|---|---|---|---|---|
| Model | $A^*$ | Vol$^*$ | $E_r$ | $T_{e,max}$ | $N_{e,max}$ |
| (1) | (2) | (3) | (4) | (5) | (6) |
| 1 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $3.8 \times 10^7$ | $4 \times 10^{12}$ | 890 |
| 2 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 3 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 4 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 5 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 6 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 7 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 8 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 9 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |
| 10 | $1.2 \times 10^{17}$ | $2.2 \times 10^{26}$ | $4 \times 10^7$ | $4 \times 10^{12}$ | 900 |

* Cross sectional areas ($A^*$), total loop volumes (Vol), and total energies ($E_r$) are calculated by assuming a loop aspect ratio $a$ ranging from 0.2 to 0.3.

* Given by $E_r = \int \pi \sigma d\tau$, $d\tau = t_1 + t_2$. Values quoted are for the case $t_2 = 0$. 

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4.1. Rapid Heating and Expansion of Coronal Gas

As soon as the flare heating is applied at the top of the loop (at time $t = 0$), the coronal gas is rapidly heated to high temperatures. Figure 4 shows that for model 8, the loop top temperature increases in about 1 s from its initial value of $10^7$ to almost $10^8$ K. The same rapid heating occurs also for models 3 and 9 (Figs. 3 and 5), although in these cases the peak temperatures are somewhat lower.

The large increase of temperature and pressure at the loop top produces a thermal wave with a conduction front that moves rapidly toward the loop footpoints. The speed of the thermal wave is much higher than the sound speed of the heated gas. For example, in model 8, the speed of the thermal wave is $2.7 \times 10^4$ km s$^{-1}$, while the sound speed is ~ 1000 km s$^{-1}$. It takes only 1.5 s for the conduction front to reach the base of the transition region. For the short loop models (models 3, 9, and 10), the conduction front reaches the loop footpoints in much less than a second. The heated coronal gas expands behind the conduction front away from the loop top at a speed of the order of only 20 km s$^{-1}$. During this initial expansion phase, the density at the loop top decreases slightly. Note that for all models, the loop is heated to its maximum temperature very rapidly to within one second from the start of the heating. This rapid heating is essentially impulsive in nature.

4.2. Heating and Evaporation of Chromospheric Material

As soon as the conduction front reaches the transition region, a large conductive flux, driven by the steep temperature gradient, is dumped into the chromosphere. Because of the inability of the chromosphere to radiate away effectively the large amount of energy deposited by the conductive flux, the
As a consequence of chromospheric evaporation, the density in the loop increases by more than one order of magnitude from its initial value. Note that once the evaporation front reaches the top of the loop, the evaporation velocity starts to drop rather quickly, even though the heating is still maintained. This occurs because the density is now much higher and further evaporation meets greater resistance. In addition, because of the increased loop density, there are more radiative losses, and less energy is available for driving chromospheric evaporation. For example, for model 8, the evaporation velocity has dropped to only 200 km s\(^{-1}\) after 30 s, even though the heating is still supplied at the maximum rate. The same occurs for the other cases. This means that the loop has its major input of evaporated chromospheric material at the very beginning of the heating phase. After the initial rapid increase of loop density caused by the initially large evaporation, the loop density increases more gradually throughout the rest of the heating phase.

4.3. Formation of Compressed Plasma at the Footpoints

As mentioned above, the localized high-pressure region generated by the conduction front acts like a two-way piston. In addition to driving the heated chromospheric material into the coronal part of the loop, the locally high pressure region also expands downward into the denser chromosphere and compresses it. The formation of compressed plasma at the base of the transition region can be seen in Figure 6 for models 3 and 8. In model 8, the compressed plasma starts to form at about \(t = 1.7\) s, when evaporation also begins. As the conduction front pushes the location of the transition region downward into deeper layers, the compressed plasma builds up, and its width increases. Initially the width of the compressed region is only \(\sim 10\) km but has doubled its size at \(t = 9\) s. The density of the compressed plasma increases by one order of magnitude from \(8 \times 10^{14}\) to \(8 \times 10^{15}\) cm\(^{-3}\). However, as the conduction front moves further downward, the density of the ambient plasma increases, and the compression effect of the advancing front becomes less effective. As Figure 6 shows, the compressed plasma can hardly be recognized at \(t = 15\) s, and it disappears entirely at \(t = 17\) s, as the ambient density becomes very high. In the compressed plasma, the temperature forms a plateau at
about \(10^4\) K. The formation of a compressed region is similar in the other model simulations.

The formation of compressed plasma at the loop footpoints was also noticed in the simulation of solar flare loops (Cheng et al. 1983; Cheng, Karpen, & Doschek 1984). Although the basic physics involved in the solar and stellar case is the same, there is an important difference. In the stellar case, because of the higher gravity, the pressure scale height is smaller than in the solar case. Therefore, when the advancing conduction front moves into deeper layers, it quickly encounters a high-density region. As a result, the formation of compressed plasma becomes less efficient than in the solar case. In the latter case, the compressed region lasts for the entire heating phase, while in the stellar case it lasts only about 10–20 s during the initial heating phase for all the models we have run. Katsova et al. (1981) suggested that the compression of plasma at the loop footpoints due to intense flare heating could be responsible for the white-light optical emission observed in many stellar flares. However, as we have seen, because of its very short lifetime and the small region involved, it is uncertain whether the compressed plasma could contribute significantly to the optical emission of stellar flares. A discussion of the various mechanisms of producing white-light emission in solar flares, such as UV irradiation and heating by electron and proton beams, and their application to stellar flares can be found in the review by Haisch et al. (1991).

4.4. Cooling Phase

In our models the flare loop is heated at a constant rate for 300 s; then the heating is stopped or decreases exponentially with a decay time constant \(\tau_d\). Figure 7 shows the temperature and density at the top of the loop as a function of time for models 3, 8, and 9, and for the different decay times. The figure shows that, as soon as the flare heating is stopped \((\tau_d = 0)\), the loop temperature drops rapidly. For model 3, the loop top temperature decreases for its maximum value of \(4 \times 10^6\) K in about 200 s, and in another 100 s it becomes as low as \(\sim 10^5\) K. This fast cooling is due to the high density in the loop \((N_e = 1.3 \times 10^{12} \text{ cm}^{-3})\), which results in very large radiative losses. In fact, the rapid cooling results in a radiative instability, as can be understood by examining the radiative cooling curve, which has a negative slope in the temperature range from \(10^7\) to about \(10^5\) K. The rapid cooling of the plasma in the loop produces large downflows with velocities of about a few hundred km s\(^{-1}\), as shown in Figures 3, 4, and 5. Because of these downflows, the loop density also drops rapidly as the figures show. In only 5 minutes, an originally hot loop becomes a cool loop.

When energy is supplied also during the decay phase, the situation is completely different. As shown in Figure 7 for the cases with \(\tau_d = 200\) and 400 s, the loop temperature decreases gradually and remains high for more than 30 minutes. The different evolutionary tracks for the various cases during the cooling phase can be best seen in Figure 8, where we plot the loop top density against the loop top temperature for models 3, 8, and 9. The arrows in the figure indicate the direction of increasing time. The figure shows that at the very beginning of the heating phase there is an immediate rise of loop top temperature from its initial value with no increase of the loop top density. The loop density starts to increase only when chromospheric evaporation begins; at that time, the loop top temperature remains approximately constant or decreases only slightly. This occurs because of the increasing radiative losses as the density increases. When the heating starts to decrease, the evolution of the loop may follow different tracks in the density-temperature plot. When there is no heating in the decay phase \((\tau_d = 0)\), the temperature and density both rapidly drop, and the loop becomes cool and dense in a very short time. This cool loop, however, is not in static equilibrium, and there are large downflows. The final state of the loop is entirely different from the initial state. On the other hand, if there is still heating in the decay phase, as in the cases for \(\tau_d = 200\) and 400 s, the loop gradually cools and eventually returns to its initial state, as shown in the figure. The reason is that the energy released in the decay phase provided the necessary radiative balance as well as pressure support for the cooling plasma thus ensuring a quasi equilibrium evolution. The two different evolutionary tracks are present also for the other models.

4.5. Effects of Changing Model Parameters

We have simulated 10 models with various loop parameters and flare energy input rates. Although the general hydrody-
namic behavior is the same for all cases, as we described above, there are quantitative differences when the loop parameters and the flare heating function are changed.

4.5.1. Changing Preflare Conditions

In Table 1, we have listed the parameters for the three initial loop models. Preflare models A and B have the same loop length ($1 \times 10^9$ cm) but different base pressures at the top of the chromosphere ($11$ and $384$ dyn cm$^{-2}$, respectively). It is clear that the higher the base pressure, the higher is the top temperature and density in the loop. Model B has a temperature of $1 \times 10^7$ K and a density of $1.3 \times 10^{11}$ cm$^{-3}$ at the top of the loop, which are about an order of magnitude greater than for model A. If the loop length is increased to $4 \times 10^9$ cm, as in model C, we obtain about the same top temperature but with a much smaller base pressure.

In order to investigate the importance of preflare conditions in the subsequent evolution of the flare, we compare model 3 and 4. These two models assume the same loop length and flare energy input but differ in the preflare conditions (they use the preflare models A and B, respectively). As shown in Table 3, the maximum temperatures and densities reached during the flare are the same for the two cases, irrespective of the initial conditions. Our numerical simulations further show that the entire flare evolution is virtually the same, except at the very beginning of the heating phase. Model 3, which starts with a lower pressure loop, gives rise to efficient chromospheric evaporation earlier and has a larger evaporation velocity, thus rapidly reaching coronal densities as high as those of model 4. The more rapid increase in density in model 3 ensures that after about 25 s the evolution of the flaring loop is about the same for both model 3 and model 4.

This behavior can be easily understood. In the initially higher density loop of model 4, radiative losses are larger and hence the flare energy deposited is largely radiated away rather than driving evaporation. Only when sufficient energy is conducted downward from the very beginning, and chromospheric evaporation can start at an earlier time.

That the hydrodynamic evolution of the flare, except during the very early stage of the flare heating, does not depend critically on the preflare conditions makes a precise knowledge of these conditions unimportant. This is a fortunate circumstance, since the physical conditions of the loop in the preflare state are poorly known, if at all, in the stellar case, where no spatial resolution is available.

4.5.2. Changing Loop Length

We have used two values of the loop length in our calculations: $2 \times 10^9$ cm in models 1, 2, 3, 4, 9, and 10, and $8 \times 10^9$ cm in models 5, 6, 7, and 8. Changing the loop length has important effects on the flare evolution. Typically, the same hydrodynamic effects occur on longer time scales in a longer loop. This can be seen by comparing the results of model 4 (short loop) and model 5 (long loop). These two models have the same initial loop temperature and nearly identical flare heating functions (see Table 2). The major effect of the loop length is on the time scales involved in the flare evolution. For model 5, it takes about 7 s for the conduction front to reach the loop footpoints, while for model 4 it takes only 0.65 s. The maximum evaporation velocities in the two models are 200 km s$^{-1}$ in model 5 and 400 km s$^{-1}$ in model 4. It takes longer to increase the density of the loop in model 5 than in model 4. The density increase at flare maximum due to chromospheric evaporation is about the same in the two cases, i.e., a tenfold increase from the initial values. The maximum temperature attained in model 5 is $6.6 \times 10^7$ K, while it is $4.3 \times 10^7$ K in model 4. Also the decay time will be longer in model 5 than in model 4 owing to reduced radiative losses (because of the lower density) and reduced conductive losses (because of the longer loop length). Thus, a flare in a longer loop will show a more gradual evolution, when everything else is about the same.

4.5.3. Changing the Heating Rate

In our simulations, we have found that the flare evolution depends significantly upon the total energy deposited during the flare but is insensitive to the way energy is deposited, if the
heating takes place at the same location, e.g., at the top of the loop. For example, the hydrodynamic response of a loop with heating deposited at a high rate (large $S_{f0}$) in a small region (small $\sigma$) is very similar to the one with small $S_{f0}$ and large $\sigma$, provided the product of $S_{f0}$ and $\sigma$ is the same for both models. The reason is that any spatial difference in the heating function will be smoothed out quickly by the large heat conductivity of the high-temperature plasma. It is only during the very initial phase of the energy release that there may be some difference.

When the flare heating rate is changed for models with identical loop length and heating spatial distributions (i.e., same $\sigma$), there are significant variations in the evaporation velocity and in the temperature and density evolution. For the short loop case, we can compare the result of models 3, 9, and 10, which all start with the same preflare loop model A and have the same temporal and spatial dependence of their heating functions, but differ in the rate of energy deposition by orders of magnitude (see Table 2). It is obvious that, as the flare heating rate is increased, the temperature of the heated plasma increases, which results in a higher temperature gradient and a faster moving conduction front. For model 9, it takes only a fraction of a second for the conduction front to reach the footpoints, while for model 10, which has a heating rate 10 times larger, it takes 5 times less. For model 3, which has a heating rate 10 times less than model 9, the time for the conduction front to reach the footpoints is, on the other hand, 5 times longer. The velocity of the chromospheric evaporation behaves similarly. The maximum evaporation speed is 900, 1700, and 2700 km s$^{-1}$ for models 3, 9, and 10, respectively. Although the rates of flare energy input for the three models are successively greater by a factor of 10, the maximum loop temperatures attained differ much less. For model 3, the maximum loop temperature is $4.3 \times 10^7$ K, for model 9 is $7.6 \times 10^7$ K, and for model 10 is $1 \times 10^8$ K.

The above results show that an order of magnitude increase of the rate of flare energy input produces only a factor of 2 or less increase in the loop top temperature. This is due to the fact that when the energy input is larger, chromospheric evaporation is also larger, thus producing a higher density in the loop. This, in turn, causes larger radiative losses thus limiting the temperature increase. In addition, when temperature increases, conductive energy transport also increases rapidly owing to its stronger temperature dependence. This again tend to limit the temperature increase. This behavior is consistent with those found previously by Rosner et al. (1978) and Jankiewicz et al. (1986). The density increase is also not as large as the flare energy input. The maximum loop top density increases by a factor of 4 from model 3 to model 9, and by another factor of 2 to model 10. In fact, it appears that the density and temperature increase becomes progressively smaller as the flare heating rate increases. The behavior of the long-loop models 5, 7, and 8, which use the same large initial loop (model C) but have different rates of energy input, is similar to that of the short-loop models just described. For example, model 8 has 10 times more energy input than model 5, yet its density is less than a factor 3 higher than in model 5. The nonlinear dependence of the density and temperature increase upon the flare energy input rate has important consequences for our interpretation of the large emission measures observed in stellar X-ray flares, as will be discussed later.

4.5.4. Changing the Decay Time Constant $\tau_d$

We have discussed previously how the evolution of a flaring loop during the cooling phase changes with time for cases with different decay time constants. When there is no energy input in the decay phase, the temperature in the flare drops rapidly, a radiative condensation instability occurs and a cool yet unstable loop forms in a very short time of about a few minutes. If, on the contrary, there is additional energy input in the decay phase, the flare loop evolves slowly, without any catastrophic instability, and eventually returns to its initial state.

5. X-RAY LIGHT CURVES AND SPECTRA

In order to compare the results from the numerical simulations with the observations and to assess how closely the theoretical loop models represent the observed properties of stellar flares, it is necessary to calculate their observational signatures. At present, most of our knowledge of the high-energy component of stellar flares comes from soft X-ray observations,
such as those from EXOSAT that we have briefly summarized in § 2 above. The EXOSAT observations were obtained with two different instruments, covering different energy bands: the Low Energy (LE), sensitive to the spectral band 0.05–2 keV, and the Medium Energy (ME) covering the range ~1–10 keV (White & Peacock 1988). In order to compare the results of our simulations with the EXOSAT observations (see Fig. 1), we have calculated LE and ME X-ray light curves as well as the predicted X-ray spectra over the energy passbands covered by EXOSAT. To this end, we have used the Raymond and Smith radiation code (Raymond & Smith 1977; Raymond 1989) for an optically thin thermal plasma in ionization equilibrium.

The Raymond-Smith radiation code calculates the X-ray spectrum in a given energy band from a plasma at a constant temperature and with unit emission measure. In our models, the temperature and the density vary along the loop. Therefore, we first calculate the X-ray spectrum from each individual computational cell along the loop using the Raymond-Smith code, and then sum up the contributions from all the cells to obtain the total X-ray spectrum from the entire loop at any given time. We can then easily obtain the X-ray light curve by integrating the spectrum over the chosen energy band. For temperatures greater than ~10^7 K, line emission from highly ionized ions of Ca and Fe makes up a large fraction of the total X-ray emission in the EXOSAT passbands. For instance, at a temperature of 4 × 10^7 K, line emission contributes about 44% of the total X-ray flux in the 1–10 keV passband, with the rest coming from continuum. The relative contribution changes with temperature; at higher temperatures, the continuum becomes more and more important.

Figure 9 shows the calculated X-ray spectra in the 1–10 keV energy band at selected times during the flare evolution for model 3 (short loop) and model 8 (large loop). As can be seen from the figure, the X-ray spectra consist of strong lines superimposed on a continuum emission. At lower energies, the emission lines are mostly from Mg xii–xii, and Fe xvii–xxiv, while at higher energies the lines are mainly from highly ionized ions, such as Ca xviii–xx, Fe xxiii–xxvi. The different spectra predicted for models 3 and 8, which have different peak temperatures, can be seen in the figure. Since the temperature of model 3 is lower than that of model 8, the contribution of lines from lower ionization stages is more important. As the temperature increases, not only the emission from highly ionized ions becomes predominant, but also the continuum emission increases greatly. The ME passband 1–10 keV contains many important diagnostic lines, notably the group of lines around 1.85 and 3.2 Å (6.7 and 3.9 keV). The group around 1.85 Å is due to the He-like Fe xxv lines and their associated dielectronic recombination lines from Fe xxiv. The group around 3.2 Å is similarly due to the lines of Ca xix and to dielectronic recombination lines of Ca xvii. The intensity ratios of the resonance lines to the dielectronic recombination lines are sensitive to the electron temperature of the emitting plasma and have extensively been used in the diagnostics of solar flares. The different line ratios within the Fe and Ca complexes for model 3 and model 8 at flare peak can be clearly seen in Figure 10. Note the strong H-like emission from Fe xxv in model 8, which has a peak flare temperature around 1 × 10^8 K. The wealth of spectroscopic information contained in the X-ray spectra will be discussed in detail in a companion paper, where we use a more refined line profile code to calculate the profiles of many diagnostically important X-ray lines including the effects of mass motions. Here we concentrate on the X-ray light curves and compare the results of our simulations with the EXOSAT broad-band LE data and with low-resolution ME data. These EXOSAT data do not have enough spectral resolution to exploit the diagnostic potential of X-ray spectra.

As we have mentioned above, line emission in the lower energy range 1–4 keV is mainly from ions of lower ionization stages than in the higher energy range 4–10 keV. This property, together with the dependence of continuum emission upon temperature, provides us with a rough method for determining the average temperature of the flare plasma using broad-band observations. Figure 11 shows the emissivity in the 1–4 keV and in the 4–10 keV bands and their ratio as a function of temperature. As the temperature increases beyond 10^7 K, the 4–10 keV emission starts to increase and around 10^8 K is dominant, primarily due to the continuum. The emissivity ratio between the two bands, as shown in the figure, is sensitive to temperature, and can be used for temperature determination over the range 10^7–10^8 K.

Figure 12 shows the light curves in the LE (0.05–2 keV) passband for models 3 and 8, and Figure 13 shows the light curves in the ME (1–10 keV) passband for models 3, 8, and 9, both for a star at a distance of 5 pc. The calculated intensities from the Raymond-Smith code and our loop models have been converted into EXOSAT count rates using the appropriate conversion factors for the LE and ME detectors. Since we are interested in the general behavior of X-ray flares, not a specific event, the use of constant conversion factors is sufficiently accurate for our purpose. Note that the conversion factor for EXOSAT ME detector does not change by more than 10% and the LE conversion factor is accurate to within 30% over.

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the relevant temperature range. The simulated light curves in unit of counts per second can thus be compared directly with EXOSAT observations (Fig. 1). In calculating the light curves, we have assumed an aspect ratio of 0.3 for our loop models. The observed aspect ratios for solar flare loops are in the range 0.2–0.5 (Cheng & Widing 1975). In the absence of spatially resolved observations of stellar flares, the aspect ratio we used in our calculations seems reasonable. We comment also that the count rates for the LE computed from our models are not as accurate as those for the ME. The reason is that the emission in the LE passband comes largely from plasma at temperatures less than $10^7$ K, which falls mostly in the flare

![Figure 10](image)

**Fig. 10.**—Detailed X-ray spectra around 1.85 Å (Fe xxiv–xxv lines) and 3.2 Å (Ca xviii–xix) calculated near flare maximum for models 3 and 8. Note the different line intensity ratios in these groups of line for the two models. The different line ratios within the group can be used as temperature diagnostics of the flare plasma.

![Figure 11](image)

**Fig. 11.**—X-ray emissivity in the (1–4) and (4–10) keV intervals, and their ratio as a function of temperature. The figure shows that the ratio can be used to determine the temperature of the emitting plasma from broad-band data.
transition region. The emission in the ME passband, on the contrary, comes from the coronal part of the loop with temperatures greater than $10^7$ K. Since the flare transition region has a very steep temperature gradient and is not spatially resolved, including a grid size as small as 10 km as used in our calculations, the emission in the LE passband often originates from only a few cells covering the transition region. As a result, the calculated LE intensity exhibits large fluctuations at any given time. The LE light curves shown in Figure 12 have been averaged over a time period of $\sim 1.5$ minutes in order to smooth out the fluctuations due to the steepness of the transition region (cf. Doschek et al. 1983). Notice that large fluctuations remain even in the time-averaged LE light curves.

We see from Figures 12 and 13 that the predicted light curves have the same rapid rise and slow decay as those observed (Fig. 1). The emission in the LE passband generally peaks later than in the ME passband, consistently with the observations. Note that a long decay in the light curve requires a slow decay in the energy input, i.e., that energy is released also in the decay phase. As soon as the energy input is stopped, the X-ray emission drops rapidly. From the comparison with observations, it appears that for most stellar flares, energy release in the decay phase is required.

At 5 pc, the count rate predicted by model 8 is high enough to account for the most intense observed stellar flares. Model 3 gives only a maximum of 0.25 counts s$^{-1}$ in the ME passband, which falls too short to explain the relatively intense flares observed by EXOSAT above its detection threshold. This is primarily due to the short loop length used in model 3, which gives a much smaller [EM] than the long-loop model 8. For model 9, the predicted ME count rate is 10 times greater than for model 3, although both models have the same loop size. The reason is that the energy input in model 9 is 10 times larger, which resulted in a factor of 3 higher density in the loop due to larger chromospheric evaporation.

Thus, even if model 9 involves only a small loop, it can explain the X-ray intensity of a medium-sized stellar flare.

6. DISCUSSION

Above, we have presented numerical simulations of the hydrodynamics of stellar flares. We have assumed, in analogy with the Sun, that stellar flares occur in loop structures. Using different loop sizes, preflare conditions, and flare energy inputs, we have calculated the hydrodynamic response of the flaring loop for 10 different cases. One important hydrodynamic result is chromospheric evaporation which results from the heating of the chromosphere by the heat flux conducted from above. The heated chromospheric material moves upward with large velocities, ranging from 400 to more than 2000 km s$^{-1}$, in the initial phase of the heating. These large velocities will result in blue-shifted components in lines of highly ionized ions, such as Fe xxiv, Ca xvii-xix, Mg xii-xiii, and others. The predicted velocities (400-2000 km) will produce blueshifts ranging from 0.002 to 0.02 Å and should be observable with the next generation of X-ray spectroscopy instruments such as those planned for AXAF. If the predicted blueshifts are indeed observed from stellar flares, this will be a strong support for the loop models we have adopted. Detailed calculations of the predicted line profiles and plasma diagnostics based on the present hydrodynamic calculations will be presented in a later paper. We note that large blueshifted Ca xix line profiles, indicating an upward mass motion velocity of about 700-800 km s$^{-1}$, have been observed in solar flares (Antonucci, Dodero, & Martin 1990).

As we have discussed above, compressed plasma is formed at the base of the transition region owing to the action of the pressure pulse which also gives rise to chromospheric evaporation. However, in contrast to the case of solar flares, this compressed plasma lasts only a few seconds during the initial phase of the flare. The main difference between the solar and stellar case is that the surface gravity on dMe stars is greater than the solar gravity. The high gravity affects the structure of the atmosphere, producing a smaller pressure scale height and, therefore, a rapid increase of density with depth. The result is that it is more difficult to form a compressed region of significant amplitude in dMe stars. It is doubtful that this region can contribute in any important way to the observed optical emission of stellar flares, in contrast to what was proposed earlier by Katsova et al. (1981). Optical flares are very intense and last up to several minutes (Petterson 1989), much longer in many cases than the lifetime of the compressed plasma in our models. We note that the problem of the origin of optical emission is still not understood even for solar flares.
in spite of the large body of observations and models now available (Mauas, Machado, & Avrett 1990).

From the predicted temporal evolution of temperature and density in the flaring loop, we have calculated X-ray light curves in the energy ranges 0.05–2 keV and 1–10 keV, which correspond to the passbands of the LE and ME experiments on EXOSAT. As we have shown above, the predicted X-ray light curves are in good agreement with the EXOSAT observations in terms of flare time profiles, X-ray count rates, and average temperatures. The X-ray observations from EXOSAT have shown that the [EM]'s of stellar flares are much larger than those of solar flares and span many orders of magnitude, from less than \(10^{51} \text{ cm}^{-3}\) to more than \(10^{53} \text{ cm}^{-3}\) (Pallavicini et al. 1990). One of the motivations of our numerical simulations was to investigate whether compact loops with sizes much smaller than the radius of a dMe star could produce the observed large X-ray [EM]'s. Figure 14 plots the predicted X-ray [EM]'s for representative loop models against the total deposited energy. Models 3 and 9 are short loops, while models 5 and 8 are large loops. The X-ray [EM]'s are those at the flare maximum and are calculated from the predicted X-ray intensities in the EXOSAT passbands. Thus the predicted [EM]'s can be directly compared with those derived from the observations. The latter are also plotted in Figure 14. The figure shows that the large-loop model 8 can reproduce [EM]'s of \(10^{53} \text{ cm}^{-3}\) as observed in the largest stellar flares. The main reason for this is primarily the large volume of the loop. Model 5, which has 10 times less energy input than model 8 produces an [EM] about an order of magnitude less, thus accounting for medium-sized stellar flares. It is not surprising to see in Figure 14 that a larger X-ray [EM] requires also a larger energy input. It is expected that a larger energy input produces more chromospheric evaporation and therefore a large loop density and [EM]. Note, however, that a larger energy input does not necessarily produce an equally larger increase of [EM]. This can be seen by comparing the small-loop models 3, 9, and 10.

As shown in the figure, the small-loop model 3 with energy input of \(10^{51}\) ergs produces an [EM] of nearly \(10^{51} \text{ cm}^{-3}\), while model 9 which has 10 times more energy input produces an [EM] one order of magnitude larger. It might be expected that another 10 times increase in the energy input could produce again an order of magnitude increase in the [EM], thus reaching nearly \(10^{53} \text{ cm}^{-3}\). The simulation shows instead that model 10 produces only a factor of 6 increase in the [EM] from model 9. In fact, as the energy input increases, the evaporation velocity also increases causing a faster increase of loop density. This in turn has the effect of hindering further evaporation of chromospheric material because of the increase radiative losses in the coronal part of the loop. Our simulations therefore clearly show that, as energy input increases, chromospheric evaporation becomes less efficient in filling the coronal loop. Moreover, even if we could increase the energy input as much as to approach the [EM]'s observed in the largest stellar flares—as we have done in model 10, there are reasons to believe that these models are not physically acceptable, as we shall see later. Our conclusion is that very large flares cannot be produced in small loops by simply increasing the energy input. Small loops instead, can produce the [EM]'s observed in small and medium size X-ray flares. This contrasts with previous conclusions by Reale et al. (1988) who argued, from the simulation of a single stellar flare, that loops of the size of a stellar radius are needed to account for the [EM] \(\sim 10^{51} \text{ cm}^{-3}\) of a small stellar flare.

In Figure 14, besides the results of our models, we have plotted the X-ray [EM]'s of stellar flares observed by EXOSAT. These data have been plotted as a function of the total X-ray energy radiated throughout the flare evolution. The close similarity that appears to exist between the theoretical and observational data points argues in favor of the physical plausibility of our loop models. Moreover, the fact that the total flare input energies are comparable, to within a factor \(\sim 2\), to the total radiative energies from the coronal part of the loop, suggests that other forms of energy, primarily that associated with chromospheric evaporation, are also comparable to the total radiative losses. Note that in our simulations we have assumed that the chromospheric radiative losses are zero, and hence we have neglected the contribution of the low-temperature material to the total radiative losses of the flare.

Figure 14 shows another interesting result. Flares with the same total energy and [EM] can be reproduced by models that differ widely in loop length (see, e.g., model 5 and model 9 in Fig. 14). Since the intensity of a flare is mainly determined by its [EM], while the time profile depends on both loop length and the unknown form of the heating function, it is not possible to determine uniquely the loop length by simply fitting an observed light curve by model calculations, as was done by Reale et al. (1988). Except possibly for the very weak and very strong flares, which require small loops and large loops respectively, the determination of the loop length will require in general additional observational constraints (such as, e.g., accurate measurements of loop density and/or evaporation velocities).

So far we have not addressed the question of magnetic confinement and realistic energy release mechanisms. In our model we have assumed that the magnetic field acts as a rigid wall to confine the heated flare plasma. Is this condition fulfilled in our models? If the plasma is to be confined by a magnetic field \(B\), the magnetic pressure \(B^2/8\pi\) must be greater than the gas pressure \(2N_e kT_e\), where \(k\) is the Boltzmann con-
magnetic fields in stellar flares, this seems unlikely. The extremely high magnetic fields required to explain large flare energies in small loops make, therefore, models such as our model 10 physically implausible.

The requirements on magnetic field strengths are much less demanding for the large loop models, primarily because of their larger volume. For the range of flare energies from $10^{33}$ to $10^{33}$ ergs, the strength of the annihilated fields ranges from 200 to 2000 G, if a volume of $3 \times 10^{-6}$ cm$^3$ is involved and we assume again that only one-tenth of the preexisting field is annihilated. Thus, it appears that for the strongest stellar flares, large loop models are needed to satisfy both the energy and confinement requirements under realistic conditions. We emphasize that the considerations above are not based on a specific model of magnetic energy release but serve only as a guide for order of magnitude estimates of the relevant physical parameters.

From the discussion above we see that our numerical simulations can explain the physical parameters and coronal emission observed from a large variety of stellar X-ray flares. Although our calculations adopt many approximations, such as the use of classical heat conductivity and a crude chromospheric model, they appear to be adequate for describing the energetics and dynamics of stellar flares. The model calculations make definite predictions about the hydrodynamic response of a stellar flare, predicting, for instance, chromospheric evaporation. The large upward velocities, on the order of several hundred km s$^{-1}$, of the heated chromospheric gas will have a definite spectroscopic signature as blueshifted components in many X-ray emission lines. These lines should be observable with future X-ray instruments which will have sufficient sensitivity and spectral resolution. The comparison between the predicted and observed line profiles will provide additional insights into the physics of stellar flares.

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