GAMMA-RAY EMISSION AND ELECTRON ACCELERATION IN SOLAR FLARES

Vahié Petrosian
Center for Space Science and Astrophysics, Stanford University, Stanford, CA 94305

James M. McTiernan
Space Sciences Laboratory, University of California, Berkeley, CA 94720

And

Holger Marschhäuser
Max-Planck-Institut für Physik und Astrophysik, Karl-Schwarzschild-Strasse 1, D-8046 Garching bei München, Germany

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ABSTRACT

Recent observations have extended the spectra of the impulsive phase of flares to the GeV range. Such high-energy photons can be produced either by electron bremsstrahlung or by decay of pions produced by accelerated protons. In this paper we investigate the effects of processes which become important at high energies. We examine the effects of synchrotron losses during the transport of electrons as they travel from the acceleration region in the corona to the gamma-ray emission sites deep in the chromosphere and photosphere, and the effects of scattering and absorption of gamma rays on their way from the photosphere to space instruments. These results are compared with the spectra from so-called electron-dominated flares, observed by GRS on the Solar Maximum Mission, which show negligible or no detectable contribution from accelerated protons. The spectra of these flares show a distinct steepening at energies below 100 keV and a rapid falloff at energies above 50 MeV. Following our earlier results based on lower energy gamma-ray flare emission we have modeled these spectra. We show that neither the radiative transfer effects, which are expected to become important at higher energies, nor the transport effects (Coulomb collisions, synchrotron losses, or magnetic field convergence) can explain such sharp spectral deviations from a simple power law. These spectral deviations from a power law are therefore attributed to the acceleration process. In a stochastic acceleration model the low-energy steepening can be attributed to Coulomb collision and the rapid high-energy steepening can result from synchrotron losses during the acceleration process.

Subject headings: acceleration of particles — radiation mechanisms: miscellaneous — Sun: flares — Sun: X-rays, gamma rays

1. INTRODUCTION

Over the years the observed spectrum of solar flare impulsive emission has been extended to higher and higher energies. The GRS instrument of the Solar Maximum Mission (SMM) observed many bursts with impulsive emission at greater than 10 MeV (Rieger et al. 1983) and few with spectra extending to 100 MeV. More recently, the SIGMA instrument aboard Granat has detected a flare up to 15 MeV (Palmel et al. 1992) and the EGRET instrument of the Compton Gamma Ray Observatory (C-GRO) and GAMMA-1 have observed gamma rays up to 2 GeV (Kocharov et al. 1991, 1994; Akimov et al. 1991; Kanbach et al. 1993). Some of these observations when combined with those at lower energies have provided very useful constraints on the impulsive phase models and on particle acceleration mechanisms.

In general, interpretation of emission at higher energies is more straightforward than at lower energies except for the complication that accelerated electrons and protons both contribute to emission at greater than 1 MeV. In the 1–7 MeV range there is contribution from nuclear line excitation, and at greater than 50 MeV there is contribution to the continuum from pion decay (see, e.g., Mandzhavidze & Ramaty 1992a). Accelerated protons are the primary source of these photons.

The observed gamma rays are the sum of these and bremsstrahlung photons emitted by relativistic electrons. The nonrelativistic counterparts of these electrons produce the hard X-rays below 1 MeV. In this range and between 7 and 50 MeV, electron bremsstrahlung is uncontaminated by emission caused by accelerated protons.

The interpretation of the bremsstrahlung emission by relativistic electrons is simple and provides constraints on the acceleration process and flare conditions. There are various reasons for this: (1) The hard X-ray and microwave emissions come from the coronal part of the loop, where an uncertain magnetic field geometry plays an important part, while the higher energy impulsive gamma-ray emissions are produced in a thin layer (a few scale height thick) deep in the photosphere; (2) the pitch angle distribution of the higher energy electrons is affected less than that of lower energy ones by Coulomb collisions; and (3) the higher energy bremsstrahlung photons are emitted primarily along the initial momentum of the electrons versus nearly isotropic emission at lower energies. Consequently, gamma-ray emission is expected to be a more true reflection of accelerated particle distribution than the hard X-ray emissions (see Petrosian 1985). We have previously analyzed such effects in detail in a series of three papers (McTiernan & Petrosian 1990a, b, 1991, hereafter MP1, MP2, and MP3, respectively). In MP3 from an analysis of center-to-limb variation of GRS gamma-ray observation (in the 300 keV to 1 MeV range).
range and the range greater than 10 MeV) we were able to set strong constraints on the pitch angle distribution of the accelerated electrons and on the field geometry.

Flares with greater than 10 MeV emission observed during cycle 21 by GRS on SMM show a strong concentration toward the limb (Rieger et al. 1983) reflecting the anisotropy of the bremsstrahlung emission at higher energies. Several unpublished reports indicate that the concentration toward the limb is not so strong for flares of cycle 22 (see, e.g., Vestrand & Forrest 1992). There are various ways that flares far away from the limb can emit a sufficiently large gamma-ray flux toward the Earth. Among these are (1) presence of a large nonradial component to the magnetic field, (2) strong field convergence or strong pitch angle scattering by plasma turbulence which will give rise to an isotropic electron pitch angle distribution and isotropic gamma-ray emission, and (3) a large contribution from pion decay, which is expected to be nearly isotropic. The two out of three recently observed flares with emission extending into the GeV range have been interpreted as pion decay emission. Kocharov et al. (1991, 1994) model the 1991 June 15 flare emission as a combination of pion decay and electron bremsstrahlung, but the 1991 March 26 flare they interpret as a pure electron bremsstrahlung event. The 1991 June 11 flare observed by C-GRO has been modeled by pion decay mechanism by Mandzhavidze & Ramaty (1992b).

In this paper we are interested in the modeling of the spectra of so-called electron-dominated flares observed by GRS (Rieger & Marschhäuser 1992; Marschhäuser, Rieger, & Kanbach 1994) which show no or very weak evidence for either nuclear line emission or pion decay emission during the two "impulsive" peaks. Almost all of the continuum emissions observed between 20 keV and 100 MeV can be attributed to electron bremsstrahlung. Presumably in such events the ratio of numbers of accelerated protons to electrons is much smaller than in other flares. Because of this absence of strong contamination by emissions due to accelerated protons, conclusions based on analysis of such events is expected to be more solidly based than for more contaminated events.

The points with error bars in Figures 1 and 2 show spectra during the two impulsive peaks from 1989 March 6 flares. (The data analysis leading to these results is described in Marschhäuser et al. 1994.) These spectra, in addition to the fact that they show a very weak sign for nuclear line emission, show significant deviations from a simple power law (see Fig. 2, in Marschhäuser et al. 1994). There is a substantial steepening of spectra toward low energies and a relatively sharp cutoff at above 50 MeV. It turns out that neither transport effects on the accelerated electrons nor effects of radiation transfer can cause deviations from power law to the extent observed if the accelerated electrons have a power-law energy spectrum. These deviations therefore must be present in the spectrum of the accelerated electron and should provide hints on the acceleration mechanism. Upturns at lower energies similar to the one seen in these figures have been observed in other flares with only hard X-ray observations (see, e.g., Lin et al. 1981; Nitta, Dennis, & Kiplinger 1990). A small fraction of such upturns in the 10–30 keV range can be attributed to reflection of bremsstrahlung photons directed toward the Sun (see, e.g., Langer & Petrosian 1977; Bai & Ramaty 1978). However, most of the excess is normally attributed to contamination due to thermal bremsstrahlung emission from a plasma heated and evaporated by the accelerated electrons (or by some other energy release process during flares). It is also possible that such

FIG. 1.—Spectrum of the first impulsive bump (13:57:29–13:58:34 UT) of the 1989 March 6 flare. The solid lines show the resultant spectrum integrated over the whole loop from a thick-target model. Electrons with a power-law energy spectrum (spectral index $\delta = 2.2$) and a Gaussian pitch angle distribution (with Gaussian dispersion $\sigma_{\theta} = 0.05$) are injected at the top of the loop with the indicated values of the coronal magnetic field $B_0$ and with a mirror ratio $b = 2.5$. These parameters are chosen to obtain the best fit for points in the 0.3–10 MeV range. The dashed lines show the optical depth effects. All these curves assume radial field lines and a heliocentric longitude of 77°.

FIG. 2.—Same as Figure 1, but for the second short impulsive spike in the interval 13:59:23–13:59:39 UT. Note that the highest energy channel point is an upper limit. Here $\delta = 2.60, \sigma^2_\theta = 0.05$, and $b = 2.5$. There are no hard X-ray data for this spike.
spectra can be the result of the acceleration process. In a stochastic acceleration model described recently by Hamilton & Petrosian (1992) it is shown that whistler waves can accelerate background thermal electrons to a few MeVs and that Coulomb collisions in the acceleration site can produce such spectral steepening at lower energies.

We shall not discuss the low-energy spectrum further here but will concentrate on the gamma-ray emissions. In § 2 we describe the effects of particle transport, including collisional and synchrotron losses, and in § 3 the effects of radiation transfer, including Compton scattering and pair production, on the gamma-ray spectrum. We show that none of these processes will result in as rapid a spectral falloff as observed at higher energies. We conclude that accelerated particles must have a rapidly declining spectrum at these energies. A possible mechanism for this is discussed in § 4. A brief summary and conclusion is given in § 5.

2. EFFECTS OF PARTICLE TRANSPORT

In order to evaluate the spectrum of bremsstrahlung radiation reaching Earth we must first evaluate the spectrum of the emitting particles. This is not the same as that of the accelerated particles because the acceleration site (presumably in the corona) is in general different from the emission site of gamma rays which is deep in the photosphere. This means we need to consider the particle transport effects. To evaluate these we need to assume a model for the accelerated electrons and flare plasma.

2.1. Model and Parameters

We assume a standard loop model for the flare such that accelerated particles are injected at the top of a symmetric loop. (For details see, e.g., Leach & Petrosian 1984; MP1.) We assume a power-law energy spectrum and a Gaussian pitch angle distribution for the injected electrons. The Gaussian width $\sigma_0$ and spectral index $\delta$ are chosen to agree with the spectrum in the middle of the observation range and to agree with previous estimates of these quantities (MP3).

The flaring loop is characterized by the magnetic field strength $B_0$ at the top and its variation along the loop specified either by $d \ln B/ds$, the logarithmic derivative of the field along the length $s$ of the loop, or by what we call the mirror ratio $(b = B_{\text{max}}/B_0)$, where $B_{\text{max}}$ is the maximum field reached by the highest energy electron considered. We assume a constant plasma density up to the transition region and then a normal solar atmospheric density profile below it (see MP2). As long as $d \ln B/ds$ is small and the column depth up to the transition region $N_{\text{eff}}$ is less than approximately $5 \times 10^{22}$ cm$^{-2}$, then gamma-ray-producing electrons will penetrate into the chromosphere and some even below the photosphere so that the details of the density and magnetic field variations become unimportant. Note that this also means that it is not important exactly where the acceleration takes place as long as it is above the chromosphere. We assume a low level of turbulence outside the accelerated region so that we can ignore pitch angle scattering of the electrons by this turbulence during the transport. As mentioned above, and as we shall see below, the high-energy electrons responsible for the observed photons with much higher than MeV energies penetrate to very high density regions below the photosphere where the density of turbulence is expected to be low. Large density of turbulence above this region will be equivalent to assuming a more isotropic pitch angle distribution for the accelerated electrons; a case which is discussed below. We will, however, discuss the role of turbulence in the process of acceleration in § 4.

2.2. Electron Transport

The transport of electrons is treated using the Fokker-Planck method. We include energy loss and pitch angle change due to Coulomb collisions and synchrotron emission which are important at low and high energies, respectively. The details of the effects of these two processes can be found in MP1. Using this method we evaluate the spectrum and pitch angle distribution of the electrons, and then, as described in MP2, we calculate the bremsstrahlung emission from these distributions as a fraction of photon energy $k$, angle $\theta$ (with respect to field direction), and distance (or column depth) along the loop. Comparison of the spectrum and directivity of the expected and observed X-ray and gamma-ray emission allowed us in MP3 to constrain the ranges of some of the model parameters mentioned above. In what follows we consider models within these bounds (see Fig. 9 of MP3).

2.3. Bremsstrahlung Spectra

We have extended these calculations to the higher energies (above 1 GeV) and have added the contribution of electron-electron bremsstrahlung. Before we compare these results with the observed spectra mentioned above we first describe some of their important features. The observations give spatially integrated spectra; therefore, we consider spectra integrated over the whole flaring loop. Figures 3 and 4 show directivity and spectra up to 5 GeV for three models. The directivity of each model is shown in Figure 3 at three energies. As expected the radiation becomes more beamed into the solar atmosphere at higher energies for all models. For a uniform field model (solid lines) most of the change in directivity occurs between 180° and

![Figure 3](image-url)
90° so that the relative directivity in the observable range between 0° and 90° shows less variation with energy. In fact the relative directivity is smaller at 2.6 GeV than at 26 MeV. This may be one of the reasons that the few observed flares with GeV emission are not all concentrated near the solar limb. An isotropic pitch angle distribution (\(\alpha_0 \rightarrow \infty\)) will give results very similar to these. Introduction of field convergence (\(b = 5\), dashed lines) in the models reduces the directivity in the backward hemisphere (toward the Sun) but increases it in the forward hemisphere. This is due to the effects of mirroring which increases the electron pitch angle and renders a more isotropic distribution in the backward hemisphere. This effect is eliminated for larger values of the magnetic field because larger pitch angle electrons quickly lose their energy by synchrotron radiation and fail to emit much bremsstrahlung gamma rays. This is demonstrated by dash dotted lines, which are for a model like the above except for an unreasonably high magnetic field \(B_0 = 10^4\) G.

The same features are also evident in Figure 4 where we plot spectra in three directions: angles \(\theta = 0°, 60°,\) and \(90°\) with respect to the (assumed) radial field lines. Spectra tend to be flat at 90° and/or for models with strong field convergence, but steepen at high energies and especially for large magnetic fields because of synchrotron losses. Note that most of these spectra do not show sharp cutoffs. The uniform field model shows considerable steepening above 500 MeV only near \(\theta = 90°\). Large steepenings are seen at lower energies only for unreasonably high magnetic fields.

Calculated spectra are compared with observations in Figures 1 and 2. The solid lines are for injected electrons with a power-law energy spectrum (spectral index \(\delta\)) and with a Gaussian pitch angle distribution (width \(\alpha_0\)). The magnetic field is assumed to converge with the specified mirror ratio. The parameters \((\delta, \alpha_0,\) and \(b)\) are chosen so that the curves give the best fit to the observed points in the 0.3–10 MeV range.

The acceptable values for field convergence and \(\alpha_0\) are within the range found in MP3.

These figures show that, if we ignore the small fluctuation due to nuclear lines in the 1–7 MeV range, the fit to the continuum in the 0.3–10 MeV range is very good for coronal magnetic field \(B_0 = 500\) G. In particular, the steepening below 1 MeV is reproduced by these models. (We note that this is not the case for all models; for example, models with uniform field, \(b = 1\), do not provide a good fit to the last feature.) Such steepenings, which are commonly observed in flares (see Vestrand et al. 1987), are due to the fact that hard X-rays and low-energy gamma rays (below 1 MeV) are emitted nearly isotropically while higher energy gamma rays are directed primarily toward the Sun. However, as is evident from Figure 1 and as shown in MP3, this cannot account for the observed steepenings in the hard X-ray region, and we must assume that electrons leaving the acceleration region must already have such a steepening in their spectra (see MP3 for details). As shown by Hamilton & Petrosian (1992) Coulomb collision by ambient plasma in the acceleration region (not those during the transport in the flare loop) can produce such spectra. Hard X-ray fluences are not available for the time interval of Figure 2.

At very high energies the electron spectrum and, therefore, the photon spectrum steepen because of the increasing influence of synchrotron losses. However, for reasonable values of the magnetic field \(B_0 \approx 500\) G this steepening is negligible at \(\approx 10^6\) MeV. An unreasonably high value of magnetic field is required for a spectral steepening as pronounced as observed above 50 MeV. The effect of synchrotron losses is to produce a gradual steepening with spectral index change of no more than 2 (see Fig. 4; MP2 eq. [40]). The energy where this transition occurs is inversely proportional to the strength of the magnetic field. This is clearly evident in Figures 1 and 2. The theoretical spectrum for \(B_0 = 5000\) G agrees with the first three data points above 10 MeV but fails to account for the low fluence observed in the highest energy channel. An even higher field \(B_0 \approx 10^4\) G is needed for this last point. These model spectra are not acceptable because they require unreasonably high magnetic fields and because they provide unacceptable fits at lower energies. The situation in Figure 2 is similar to that in Figure 1, except that here the results are not as clear-cut because the highest energy datum is only an upper limit and because the spectrum around 10 MeV is noisier. We conclude that synchrotron losses during the transport cannot be responsible for the rapid fall of spectra at high energies.

More isotropic distribution of accelerated electrons, stronger convergence of the field, or presence of pitch angle scattering by turbulence could have an effect similar to that of higher magnetic fields. Here again, however, the steepening of spectra at high energies will be gradual and not as abrupt as observed. We therefore conclude that, in general, transport effects cannot account for the rapid spectral steepening at high energies.

3. EFFECT OF RADIATION TRANSFER

The higher energy electrons penetrate deeper into the solar atmosphere so that their bremsstrahlung emission must go through a larger column depth in order to reach Earth. Electrons of energy \(E\) (in units of \(m_c^2\)) and pitch angle \(\mu\) lose most of their energy via collision by the time they reach a column depth (Leach & Petrosian 1981) of

\[
N(E) = (4\pi r_0^2 \ln \Lambda)^{-1} \mu E^2/(E + 1),
\]  

(1)
where $r_0 = 2.8 \times 10^{-13}$ cm and $\ln \Lambda \approx 20$ is the Coulomb logarithm.

If the magnetic field strength is large,

$$B > 2500 \text{ G(n/10^{14} \text{ cm}^{-3})(10^3/E)} ,$$

then synchrotron losses become important and the column depth penetrated will be less than that given in equation (1). The column depth penetrated could be even smaller than this if the electrons are scattered (e.g., by turbulence) or if the field convergence is large. Nevertheless, the highest energy particles ($E \approx 10^9$) with $\mu = 1$ will penetrate to a column depth $N \approx 10^{25}$ cm$^{-2}$. This is larger than the inverse of the Thompson cross section $\sigma_0^{-1} = 1.5 \times 10^{14}$ cm$^{-2}$ so that the effects of scattering or absorption of photons emitted by such electrons must be investigated. In fact, since $N(E)\sigma_0 \approx (2/3 \ln \Lambda)E^2/(E + 1)$, then for $\ln \Lambda = 20$ radiative transfer effects may become important at energies of 15 MeV or higher. For investigation of these effects we need to evaluate the emission, absorption, and scattering coefficients.

In what follows we shall first use the approximation that most of the photons of energy $k$ originate from the column depth where electrons of the same energy lose most of their energy. For radial magnetic field lines this means that photons of energy $k$ are emitted from a column depth approximately as given by equation (1) with $E$ replaced by $k$. For nonradial field lines, the photon column depth will be smaller by a factor cos $\eta$, where $\eta$ is the angle of field lines with respect to the radial direction.

### 3.1. Estimated Optical Depths

Two processes will dominate the scattering and absorption of high energy photons.

**Compton scattering** will dominate at low energies with cross section $\sigma \sim \sigma_0$ and with initial and final photon energies comparable, $k \sim k'$. As we approach relativistic energies the cross section is given by the Klein-Nishina formula, which is comparable to $\sigma_0$ only for scattering angles no more than $(2k)^{-1/2}$ (with $k \approx k'$), but decreases as $\sigma_0/k$ for larger angles (with $k \gg k'$) in analogy with the forward scattering is elastic it will not attenuate the emitted bremsstrahlung photon flux. Only large-angle scattering attenuates the flux but with an effective optical depth of

$$\tau_0(k) = N(k)\sigma(k) \approx (2/3 \ln \Lambda)\mu \approx 0.03\mu \ll 1.$$  

Or if instead we use the total Klein-Nishina cross section, $\tilde{\sigma} \approx (3\sigma_0/8k)\ln 2k$, we find that for $\mu = 1$ the optical depth $\tau(k) > 1$ for $k \geq \Lambda^2/2$, which is much larger than energies under consideration here as long as $\ln \Lambda > 2$. The variation of $\ln \Lambda$ with $E$ based on formulae presented in Leach & Petrosian (1982) for a neutral gas with solar abundance is shown in Figure 5. From this we conclude that scattering effects will be unimportant, and even when important will be nearly independent of energy and not capable of producing a sharp spectral cutoff.

**Pair production** in the electronic and nuclear Coulomb field will be the most important process at energies above 1 MeV (see, e.g., Hubbel 1977). The Bethe-Heitler Born approximation cross section for pair production varies approximately as $\sigma_0 \propto (\ln k - 2)$. Figure 5 shows a plot of the total cross section (mostly due to pair production) from 0.1 MeV to 100 GeV for a plasma with photospheric abundance calculated from compilations by Hubbel, Gimm, & Overbo (1980). As is evident in the range from 10 MeV to 10 GeV of interest here, the cross section is nearly constant and can be approximated by $\sigma \approx 0.04\sigma_0$. Consequently the effective optical depth at high energies ($k > 20$) is given by

$$\tau_0(k) \approx 1.3 \times 10^{-3}(20/\ln \Lambda)\mu k ,$$

which exceeds unity above 400 MeV for $\mu = 1$.

The above optical depths are for photons radially moving out of the photosphere, which is the case when the flare is at the disk center. Most of the high-energy gamma-ray-producing flares are observed to occur away from the center of the solar disk. The flares of Figures 1 and 2 took place at heliocentric longitude $\phi = 77^\circ$. In general for a longitude $\phi$ the optical depth is given by $\tau_0(k) = \tau_0(k)/\cos \phi$, where $\tau_0$ is the radial optical depth discussed above. This equation is valid for $\cos \phi > (2h_\perp R_\odot)^{1/2}$, where $h_\perp$ is the density scale height at the emission site. For $h_\perp < 10^3$ km this expression is valid for flares to within 4° of the limb. The optical depth at the limb will be about $\tau_0(k) \approx (R_\odot/2h_\perp)^{1/2} \tau_0(k) \approx 17 \tau_0(k)$ and will exceed unity above 5 MeV. At $\phi = 77^\circ$ the effective Compton optical depth (eq. [3]) is still negligible but the effective pair production depth (eq. [4]) for $\mu = 1$ is $\tau_0(k) \approx 0.006k$ and will exceed unity at photon energies greater than 100 MeV. From this we conclude that a careful analysis of the radiative transfer effect is required at energies above tens of MeV.

### 3.2. Spatial Variation of the Source Function

For a detailed analysis of the spectrum emerging from a thick-target optically thick bremsstrahlung source we need to consider the angular and spatial distribution of the emitted photons. Here we need the variation along the loop of the bremsstrahlung emissivity $J(k, \theta, \Lambda)$, where $J(k, \theta, \Lambda)d\Lambda$ is the emissivity integrated over the cross section of the loop from column depths $N$ to $N + dN$ at an angle $\theta$ with respect to the...
magnetic field lines, which are assumed to be radial for high-
energy emission under consideration here. In Figure 6 we plot
the quantity $NJ(k, \theta = 0, N)$ versus $N$ (in units of effective
column depth of eq. [1]). In a log-log plot this demonstrates
where most of the photons are generated more clearly than
plots of $J$ versus $N$. For the uniform field model (Fig. 6a)
the spatial distribution of the emission is a simple function
of the column depth at $\theta = 0^\circ$ and the shape of this distribution
is fairly insensitive to the photon energy $k$. The quantity
$NJ$ rises nearly linearly at low $N$ until it reaches a maximum at
$N_{\text{max}}(k, \theta)$ after which it drops fairly rapidly. This behavior
can be approximated by the following simple analytic expression

$$J(k, \theta, N) = J_0(k, \theta)[1 + N/N_{\text{eff}}(k, \theta)]^{-(v/N_{\text{eff}})}, \tag{5}$$

where $J_0(k, \theta)$ is the emissivity from the whole loop in the
optically thin regime (Figs. 1 and 2, solid lines).

For hard X-rays ($k \ll 1$) there is little anisotropy. For the
angularly averaged quantities at these energies, Leach (1984;
see also Canfield et al. 1986) finds empirically that the above
expression with $v \geq \delta/2$ and $N_{\text{eff}}(k)$ given by equation (1)
provides a good approximation for the numerical results. Equation
(18) in MP1 gives electron spectra integrated over pitch angles as a function of dimensionless depth $N/N_0$, [$N_0 = 5 
\times 10^{22} \text{ cm}^{-1}(20/\ln \Lambda)$]. Integrating the nonrelativistic expres-

![Figure 6a](image)

**Fig. 6a**—(a) Variations with column depth (in units of the energy-dependent effective depth of eq. [1]) of the bremsstrahlung emissivity per column depth $J(k, N, \theta)$ for the uniform field model with $\delta = 3, B_0 = 500 \text{ G}$, and $\alpha_0^2 = 0.60$, for photon energies 26 MeV and 2.6 GeV, for angles $\theta = 0^\circ$ and $90^\circ$. The diamonds show the same for the approximate expression in eq. (5), for $v = 2$ and $N_{\text{eff}} = (v - 1)N_{\text{max}}$. To show where most of the emission occurs in this log-log plot we have plotted $NJ$ instead of $J$. Note also that for high-energy photons $N_{\text{eff}}$ is larger, so that they penetrate to larger column depths. (b) Same as (a), but for the converging field model with $\alpha_0^2 = 0.01, \delta = 3, b = 5.0$, and $B_0 = 500 \text{ G}$.
3.3. Radiation Transfer

For relativistic energies we can approximate the radiation transfer by equation for a semi-infinite plane-parallel atmosphere extending from the optical depth zero (at the observer) to infinity (in the core of the Sun). We assume a semi-infinite cylindrical source confined to a footprint of a magnetic loop with the field lines having an average angle \( \eta \) with respect to the local vertical (or direction of stratification). Then, when viewed at an angle \( \theta \) with respect to the vertical (except for \( \theta \approx 90^\circ \)), we can define an absorption optical depth

\[
d\tau(k) = \sigma(k) \cos \eta \cos \theta \, dN,
\]

where \( \sigma(k) \) is the cross section for pair production and \( dN \) is the increment of column depth along the field lines. The radiative transfer equation integrated over the cross section of the magnetic loop can then be written as

\[
dl(k, \theta, \tau) = -J(k, \theta, \tau) + \frac{J(k, \theta, N)}{\sigma(k)} \left(1 - e^{-\tau_{\text{eff}}(k)}\right),
\]

where \( J(k, \theta, N) \) is the emissivity per column depth integrated over the loop cross section. The solution of this is simple and yields the radiation intensity at \( \tau = 0 \):

\[
I(k, \theta) = \int_0^\infty \frac{J(k, \theta, N)}{\sigma(k)} e^{-\tau} \, d\tau.
\]

Assuming bremsstrahlung as the main source of emissivity, the variation with depth of the emissivity can be approximated by equation (5). As mentioned above, this approximation will be sufficient for our purpose here. It will deviate from the results obtained from more exact numerical emissivity profiles shown in Figure 6 by less than the observational uncertainties. With this approximation then, it can be shown that

\[
I(k, \theta) = J_\nu(k, \theta) e^{\tau_{\text{eff}}(k, \theta)} E_\nu(\tau_{\text{eff}}(k, \theta)),
\]

where \( E_\nu(x) \) is the exponential integral function of order \( \nu \).

As expected, at small and large optical depths \( I(k) \approx J_\nu(1 - \tau_{\text{eff}}) \) and \( I(k) \approx J_\nu/\tau_{\text{eff}}, \) respectively. For intermediate depths \( I(k) \approx J_\nu/\tau_{\text{eff}} + \nu - 1 \) is a good approximation. This demonstrates that the most important parameter determining the shape of the observed photon spectrum is the energy dependence of the effective optical depth. As discussed above, for Thompson or Compton scattering \( \tau_{\text{eff}} \) is nearly constant, but for pair production \( \tau_{\text{eff}} \propto k \). The dashed line in Figure 7 shows variation with photon energy of \( \tau_{\text{eff}} \) and the ratio \( I(k)/J_\nu(k) \) assuming \( \theta = 77^\circ \) and \( \nu = 3 \). The difference between such curves for different \( \nu \)-values is well below the accuracy demanded by the data, and the most important parameter in determining the shapes of the spectra is the energy and angular dependence of \( \tau_{\text{eff}}(k, \theta) \).

As expected, the spectrum steepens with an increase of one unit in the photon spectral index at higher energies. At first glance this kind of break seems promising, but as shown by the dashed lines in Figures 1 and 2 this yields a gradual spectral steepening which is not in agreement with the rapid falloff observed for this flare. One way optical depth effects can produce a sharper cutoff would be if the emissivity were a narrow function of the column (or optical) depth. For example, for \( J(N, k) \propto e^{-\delta[k - \tau_{\text{eff}}(k)]} \) then \( I(k) \propto e^{-\tau_{\text{eff}}(k)} e^{-e^{-\tau_{\text{eff}}(k)}} \). However, this is not the case for the thick-target model for solar flares. It is true that the emissivity especially at high
energies is confined to a thin layer with a width of a few density
scale height $h \ll R_\odot$. However, as shown in Figure 6, the
column depth variation over this distance is significant.

We should also mention the fate of the pairs of electrons and
positrons produced in the absorption process. Some small fraction
of their energy will go into production of secondary bremsstrahlung photons of lower energy than the energy of the primary photons. Since the cross section is larger at lower energies, the secondary photons will also be absorbed and produce even lower energy pairs. However, the bulk of the energy of the pairs will be lost by collisions near where they are produced or in the other footpoint of a closed magnetic field
loop. Thus very little energy of absorbed photons will appear as
observable lower energy photons.

We therefore conclude that radiation transfer effect cannot
provide the steep cutoff observed in this flare.

4. ACCELERATION MODEL

The deviations of observed bremsstrahlung spectra from
model calculations described above signify breakdown of some
of the assumptions made in the calculations. Among these the
most likely candidate is the assumption of a pure power law for the
energy spectrum of the flux of injected (or accelerated)
electrons. We have evaluated the effect of other assumptions,
like the pitch angle distribution, the magnetic field geometry and
its variation, plasma density profile, and so forth, and find
that none of these can account for the above deviations as long
as the parameters describing them are within a reasonable range,
in particular, if they are within the ranges determined by
previous observations (e.g., those in MP3).

One possibility is that the acceleration process is slow com-
pared to the duration of the impulsive phase and there is not
sufficient time to accelerate particles to energies much greater
than a few tens of MeV. This may be a possible explanation for
the second impulsive burst shown in Figure 2 which lasted not
much more than a few seconds. However, the first impulsive
burst (Fig. 1) lasted more than a minute, in which case, under
this hypothesis, one would expect a gradual extension of the
spectrum to higher and higher energies as the burst progressed.
This is not what is seen. The time profile of different channels
are very similar, indicating that the observed deviations from a
power law are a result of the details of the acceleration process.
Thus, we need to specify this process.

In this paper we consider stochastic acceleration by plasma
turbulence (whistlers and Alfvén waves) which in its simplest
form results in a power-law spectrum of particles. However,
this involves numerous simplifying assumptions. Deviations
such as those discussed above could arise in models with more
realistic assumptions. One possibility is to relax the assumed
energy dependence of acceleration rate relative to that of
escape time (Petrosian 1994). Another is to include the effects
of competing processes. Two processes that can be important
in a flare plasma are the Coulomb collisions and synchrotron
losses. As discussed in § 2 these are the primary energy loss
mechanisms during the transport of the electrons. These pro-
cesses could play important roles during the acceleration as
well. We now consider the effects of these processes in the
acceleration site as the cause of the deviations from a power
law of the initial acceleration.

As usual the effect of collisions is felt at lower energies, and
as mentioned above, this can easily explain the softening of
the accelerated electron spectra at low energies which is required
for production of the observed hard X-ray excesses. The devi-
ation at low energies from the power law occurs at energies $E < E^*_\gamma$ where the collisional loss rate is comparable to the
acceleration rate. For further details see Hamilton & Petrosian
(1992). Since the synchrotron loss rate increases quadratically
with energy its effects will be felt at higher energies, and we
expect that at sufficiently high energies, $E > E^*_\gamma$, the synchro-
ton loss rate will be larger than the acceleration rate. This will
cause an exponential decrease in the number of electrons at
higher energies. Similar explanations, with some analytic
expression describing both of these effects, are also given by

If the stochastic acceleration is the correct mechanism, then
it must accelerate electrons within a timescale which is equal to
or shorter than the timescale of modulation of the observed
gamma rays. The rise times of gamma rays in Figures 1 and 2
are about 30 s and no more than 2 s, respectively. Furthermore,
if synchrotron loss in the acceleration site is to be responsible
for the high-energy spectral cutoff at about 50 MeV, then its
timescale must also be comparable to these timescales. For $E^*_\gamma = 10^2$ MeV the 2 s timescale spike of Figure 2 requires a
magnetic field of about $10^3$ G. This is somewhat large if the
acceleration is to occur in the corona. A lower, more reason-
able ($\sim 300$ G) magnetic field will be required for the slower
spike of Figure 1. Alternatively, when the acceleration process
is slower (e.g., when density of plasma turbulence is low), the
high-energy spectral steepening will be more pronounced and
will occur at a lower energy. Consistent with this, the flux in
the last two channels in Figure 1 falls off more steeply than in
Figure 2.

The detailed and more quantitative comparison of this
picture with observations is beyond the scope of this paper.
Our aim here has been to show that electron transport and
radiative transfer effects cannot be the source of rapid spectral
steepening at high photon energies. In a separate publication
we shall consider the effects of synchrotron losses on the spec-
trum of the accelerated electrons under various plasma condi-
tions.

5. SUMMARY AND CONCLUSIONS

We have extended our earlier calculation of bremsstrahlung
emission by accelerated electrons to greater than GeV photon
energies. This work is motivated by recent detection of impul-
sive flare emission up to 2 GeV by various instruments. We
have concentrated on the comparison of the model spectra
from the so-called electron-dominated flares observed by the
GRS instrument on SMM. These flares show no or a negligible
sign of nuclear line emission, indicating smaller numbers of
accelerated protons as compared to electrons. Consequently,
the spectrum at high energies is expected to be uncontami-
nated by pion-decay gamma rays and must be solely due to
bremsstrahlung by relativistic electrons. These spectra show
distinct softening at low energies (as in many other flares) and
a rapid falloff at the highest energy channels.

We have shown that in the standard thick-target model,
electron transport effects cannot account for these features,
even for unreasonably high values of magnetic field, if the
accelerated electrons have a power-law spectrum. Some spec-
tral steepening (or flattening) at certain viewing angles is
expected but not to the degree observed here. In general, the
transport effects, those considered here and others (e.g., strong-
er field convergence, scattering by turbulence, etc.), can
produce only gradual steepenings which are less pronounced
than that observed, for example, in the 1989 March 6 flare.
We next consider the radiative transfer effects (scattering and absorption of bremsstrahlung protons) which become more important at higher energies because such photons are emitted by higher energy electrons which penetrate larger column depths into the photosphere. We have shown that Compton scattering is unimportant but pair production absorption can be important especially for flares near the limb. However, we show that the effects of absorption produce a gradual steepening over a wider band (changing spectral index by one unit) which is not as sharp as observed.

We therefore attribute the spectral deviations from a pure power law to the acceleration mechanisms. The low-energy steepening is similar to that observed in other flares and, as shown by Hamilton & Petrovian (1992), can be caused by Coulomb collisions in the acceleration site (not collisions during the transport from the acceleration region in the corona to radiation regions, which tend to be the higher density chromosphere and photosphere). We have considered the possibility of synchrotron losses (again in the acceleration site not during the transport) as a cause of steep cutoff in the spectrum of accelerated electrons at high energies. We show that this requires a high value of magnetic field for spectral cutoff above 100 MeV and for short duration (less than 1 s) bursts. A more detailed analysis of this problem will be given elsewhere.

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