THICK TARGET BEAM INTERPRETATION OF STEREO OBSERVATIONS OF A
SOLAR HARD-X-RAY BURST

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ABSTRACT

Observations by Kane et al. of an occulted solar hard-X-ray burst by two separated spacecraft
are analyzed in terms of the thick target beam model. It is pointed out that the much higher flux
of low-altitude emission compared to that from high altitudes is consistent with the thick target
model for downward injection of electrons from a region lying at an ambient column number

\[ \Delta N = 2.5 \times 10^{18} \text{ cm}^{-2} \]

above the \( z = 25,000 \) km level. The actual electron injection rate required
above 10 keV is \( \dot{N}_e \approx 1.5 \times 10^{38} \text{ s}^{-1} \). For the associated return current to be stable against ion-

sound wave generation, the beam area \( A \), ambient density \( n \), and temperature \( T \) must satisfy

\[ \frac{AnT^{1/2}}{112} > 10^{32} \text{ cm}^{-1} \text{ K}^{1/2} \]

and in particular the acceleration site density must far exceed \( 10^9 \text{ cm}^{-3} \).

From simultaneous (occulted) soft-X-ray data, we find that the parameters lie around this limit, sug-
gesting that return-current instability may be acting as an electron flux limiter.

Subject headings: plasmas — Sun: X-rays — X-rays: bursts

I. INTRODUCTION

The interpretation of solar hard-X-ray bursts has important implications for overall flare energetics and particu-
larly for the efficiency required of the acceleration mechanism (for the most recent reviews see Kane et al. 1980b;
Brown and Smith 1980). Furthermore, the implications are strongly dependent on the interpretive model used, but
at present there is insufficient data for an unambiguous distinction between the three models prevalent in the litera-
ture. These are (a) the thick target beam model (e.g., Brown 1971, 1972; Petrosian 1973; Lin and Hudson 1976);
(b) the trap-plus-precipitation model (Hudson 1972; Kane 1974; Melrose and Brown 1976); and (c) thermal models,
either dissipative (Brown, Melrose, and Spicer 1979; Smith and Lilliequist 1979; Smith and Auer 1980) or adiabatic
(Mätzler et al. 1978). This continuing ambiguity of interpretation can be partly attributed to the lack of spatial
resolution of hard-X-ray burst data, a problem currently being eliminated with the recent advent of hard-X-ray
imaging. In the meantime, however, Kane et al. (1979) obtained the first observation ever of the distributed height
structure of a burst by means of two detectors separated in solar longitude—Pioneer Venus Orbiter (PVO) and
International Sun-Earth Explorer-3 (ISEE). This represents a great advance on previous data on height structure,
which was limited to observations by single spacecraft of the emission from the single height range above the occult-
ing solar limb (e.g., Hudson 1978).

The burst observed (UT 0631 on 1978 October 5) was estimated by Kane et al. (1979) to come from an active
region 15° beyond the solar limb. Consequently ISEE, which lay near the Sun-Earth line, could observe only emis-
sions from heights \( z > z_1 = 25,000 \) km above the photosphere. PVO, on the other hand, lying 12°5 to the east of the
Sun-Earth line, could observe all \( z > z_2 \approx 700 \) km. The essence of these data can be summarized as follows (cf.
Kane et al. 1979, Fig. 1b):

A. The ISEE flux had a power-law spectrum,

\[ \left( \frac{dJ}{de} \right)_{\text{ISEE}} = 2.1 \times 10^{4} \text{ cm}^{-2} \text{ s}^{-1} \text{ keV}^{-1} \]

over the observed band from 5 to 50 keV.

B. The total PVO flux in the observed range 50–500 keV was

\[ \mathcal{J}(>50 \text{ keV})_{\text{PVO}} \approx 50 \pm 1 \text{ cm}^{-2} \text{ s}^{-1} \]

or about 600 times the corresponding ISEE flux obtained from equation (1) by extrapolation above 50 keV (Kane
et al. 1979). This ratio may be in error by a systematic factor of about 2 for instrumental reasons.

C. The PVO spectral information is less definitive because of the small number (three) and very broad energy
ranges (50–100, 100–250, 250–500 keV) of the channels, and because of variability in the collecting area and un-
certainly in energy calibration. If we take the three measured $dJ/d\varepsilon$ values with errors purely due to count statistics (S. R. Kane, personal communication), we find that no single power law can fit adequately, since the spectrum steepens above 150 keV. This may be attributable to a genuine steepening of the electron spectrum at high energies or to occultation of the emission by the highest energy electrons even as seen by PVO (cf. § II). On the other hand, if we take the deviations of the three points from the best-fit regression line as a measure of the errors, we obtain $(dJ/d\varepsilon)_{PVO} \approx 10^{4.7 \pm 1.4} \varepsilon^{-2.6 \pm 0.7}$. Thus using single power-law fits, the most we can say is that the PVO spectrum is harder than the ISEE spectrum by about 2 in the spectral index and the total PVO flux is about 600 times the ISEE flux, but with uncertainty possibly up to one order of magnitude. For the moment we use a (round figure) spectrum $(dJ/d\varepsilon)_{PVO} \approx 10^{-5} \varepsilon^{-3}$ to represent the PVO data and discuss the effects of uncertainties again later.

Kane et al. (1979) mention two possible interpretations of the large flux discrepancy in (B), namely that the emission is highly anisotropic or that it is concentrated at altitudes <25,000 km and so is mostly occulted from observation by ISEE. They discount the former interpretation on the grounds that the required directivity is improbably high and is inconsistent with statistical data on center-to-limb burst variations (e.g., Datlowe et al. 1977; Kane et al. 1980a). The argument of improbability is in fact seen to be conclusive when one examines available theoretical work on directivity. Even for a perfectly collimated monoenergetic beam (cf. Koch and Motz 1959), a bremsstrahlung intensity variation of 600 times over 125° could occur only for precise alignment of the beam opposite the line of sight. For a power-law spectrum the maximal computed flux variations at 50 keV even over 90° for a downward directed beam are only a factor of about 5 when beam scattering is neglected and about 3 when scattering is included (Brown 1972; Haug 1972; cf. Langer and Petrosian 1977; Bai and Ramaty 1978).

Considering the explanation of (B) in terms of occultation, Kane et al. (1979) give a broad interpretation as follows. They suppose that the ISEE flux arises by continuous injection of electrons into a low-density ($n \text{ cm}^{-3}$) coronal region of thin target bremsstrahlung where they have a lifetime $\tau$. Then, on the assumption that electrons are injected at the same rate into the dense chromosphere to produce thick target bremsstrahlung observed by PVO, they conclude from the flux ratio in (B) that $n \tau \approx 10^5 \text{ cm}^{-2} \text{s}$. Furthermore, from the independence of burst decay time on photon energy, they conclude that $\tau$ is probably governed by “escape” from the thin target region, rather than by collisions.

In this Letter we discuss these general results quantitatively in terms of a specific model situation, namely the thick target beam model mentioned above (a) (cf. Brown 1980). In this model electrons are accelerated from some height $z_0$ with small pitch angles and descend into the chromosphere to emit thick target bremsstrahlung there, emitting thin target bremsstrahlung as they pass through the low density corona. Here electron “escape” is thus downward to the chromosphere rather than away from the Sun. This model automatically satisfies the assumptions made by Kane et al. (1979), that the electron fluxes into both regions are one and the same, with $\tau$ = transit time across the thin target region (i.e., between $z_0$ and the lowest height $z_1$ observed by ISEE). We use the results of Brown and McClymont (1975) for the theoretical distribution of flux with height in the thick target model.

II. THICK TARGET INTERPRETATION

First, we assume that electron energy losses are collision dominated (cf. a posteriori check) and note (cf. Brown and McClymont 1975, Table 1) that the typical mean “coronal column density” in the height range seen by ISEE is $N_1 = \int_{z_0}^{z_1} ndz < 10^{19} \text{ cm}^{-2}$ even for active region densities. This has negligible effect on electrons of energy $\geq 5 \text{ keV}$, so the region can be rigorously regarded as thin target for the entire range of emission seen by ISEE.

Second, the column density seen by PVO is

$$N_2 = \int_{z_0}^{z_1} ndz \gtrsim 5 \times 10^{21} \text{ cm}^{-2}$$

even for quiet-Sun densities, which is enough to completely stop electrons of up to about 200 keV. Thus the region observed by PVO is thick target for the spectral range covered by at least the first two channels of PVO, though some occultation may occur at third channel energies (cf. § II discussion of spectrum).

We emphasize of course that the whole concept of a unique column density is strictly valid only in plane-parallel structures in contrast to the strong inhomogeneity of flares. Thus our parameter values and conclusions apply only in some mean sense similar to that in flare atmosphere models.

Suppose then that electrons are injected from $z_0$ at a rate

$$\bar{\Phi}(E_0) = (\delta - 1) \frac{\Phi_1(E_0)}{E_1} E_0^{-4} \text{s}^{-1} \text{per unit } E_0,$$  

(3)

where $\Phi_1$ is the total injection rate for $E_0 \geq E_1$. Then with $\delta = 4$, we multiply equation (8) of Brown and McClymont (1975) by the area $A$ of injection to find the predicted thick target flux observed by PVO, viz.,

$$\left( \frac{dJ}{d\varepsilon} \right)_{PVO} = 7 \times 10^{-32} \Phi_1 \varepsilon^{-2} \text{ cm}^{-2} \text{s}^{-1} \text{ keV}^{-1}.$$  

(4)
where \( \Phi \) is \( \Phi (E_1 = 10 \text{ keV}) \) and \( e \) is in keV. Hence, theoretically

\[ [J(e \geq 50 \text{ keV})]_{I S E E} = 1.42 \times 10^{-2} \Phi_0 \text{ cm}^{-2} \text{s}^{-1} \text{ keV}^{-1} \].

(5)

Again, multiplication of equation (9a) of Brown and McClymont (1975) by equation (4) and by the injection area gives the thin target emission per unit height of the target, viz.,

\[ \left( \frac{d^2J}{dz} \right)_{\text{thin}} = 2.4 \times 10^{-6} \Phi_0 e^{-e} n(z) \text{ cm}^{-2} \text{s}^{-1} \text{ keV}^{-1} \].

per cm column,

(6)

so that the ISEE (total thin target) flux should be

\[ \left( \frac{dJ}{dz} \right)_{I S E E} = \int_{z_1}^{z_2} \left( \frac{d^2J}{dz} \right)_{\text{thin}} dz = 2.4 \times 10^{-6} \Phi_0 \Delta N e^{-e} \text{ cm}^{-2} \text{s}^{-1} \text{ keV}^{-1} \],

(7)

where \( \Delta N \) is the total thin target column density in the ISEE field. Thus, theoretically,

\[ [J(e \geq 50 \text{ keV})]_{I S E E} = 9.6 \times 10^{-22} \Phi_0 \Delta N \text{ cm}^{-2} \text{s}^{-1} \].

(8)

From theoretical equations (4)–(8), we have the following: The ISEE spectrum (7) should be two powers steeper than the PVO spectrum (4), owing to the collisional hardening of the electron flux in the thick target region (cf. Brown 1971; Brown and McClymont 1975, especially their Fig. 5), which is consistent with the observed spectra as described in (C).

According to equations (5) and (8), the flux ratio observed should be

\[ \frac{[J(e \geq 50 \text{ keV})]_{I S E E}}{[J(e \geq 50 \text{ keV})]_{P V O}} = 6.8 \times 10^{-22} \Delta N \],

(9)

which is in agreement with the observations (B) and (C) if a value

\[ \Delta N = 2.5 \times 10^{18} \text{ cm}^{-2} \]

(10)

is assigned to the plasma column density between \( z_1 = 25,000 \text{ km} \) and the height \( z_0 \) of the acceleration (injection) site. Setting \( n = \text{mean density in this height range} \{ z_0, z_1 \} \) and \( \Delta z = z_0 - z_1 \), we can see the correspondence between our particular result (10) and the Kane et al. (1979) general conclusion that \( n \tau \approx 2 \times 10^8 \text{ cm}^2 \text{s}^{-1} \) in the thin target region, as follows. In our model the mean electron (escape) lifetime is \( \tau = \Delta z / v = \Delta N / n v \), where \( v \) is the mean electron speed, \( \sim 1.3 \times 10^{10} \text{ cm} \text{s}^{-1} \) at 50 keV. Thus we find \( n \tau = \Delta N / v = 2 \times 10^8 \text{ cm}^2 \text{s}^{-1} \) also.

The value (10) for \( \Delta N \) provides a valuable constraint on conditions at the electron source. First, if we assume a density \( n(=10^{19} n_0) \), we obtain the height of the acceleration region, viz.,

\[ z_0 = z_1 + \Delta z = \left( 1 + \frac{1}{n_0} \right) \times 25,000 \text{ km} \].

(11)

For quiet-Sun conditions at these heights (cf. Brown and McClymont 1975), the typical density \( n_0 \approx 0.1 \), which would imply \( z_0 = 275,000 \text{ km} \). However, active region densities are typically one order higher, i.e., \( n_0 \approx 1 \), for which the acceleration region would be located at \( z_0 = 50,000 \text{ km} \) or \( \Delta z = 25,000 \text{ km} \) above the minimum altitude seen by ISEE. Second, however, we may further constrain \( n \) and hence \( z_0 \) both from theoretical considerations and from independent data, as follows.

First, from equations (2) and (4) we find the total electron injection rate above 10 keV to be

\[ \Phi_0 \approx 1.5 \times 10^{18} \text{ s}^{-1} \],

(12)

with a corresponding electron flux,

\[ F_0 = \Phi_0 / A = 1.5 \times 10^{17} / A_{19} \text{ cm}^{-2} \text{s}^{-1} \],

(13)

where \( A (=10^{19} A_{19}) \) is the cross-sectional area of the beam. In order for this flux to be sustained as a beam, it is necessary that the neutralizing return current should not become microscopically unstable (see, e.g., Brown and Melrose 1977; Hoyng, Knight, and Spicer 1978) to generation of ion-acoustic or electrostatic ion-cyclotron waves. The stability condition for the ion-acoustic instability depends on the ratio of the return current drift velocity \( \nu_r \) to the ion-sound speed \( \nu_s = (kT_i/m_i)^{1/2} \), and on the ratio of electron to ion-temperature, \( T_e/T_i \), in the background plasma according to (for stability):

\[ \nu_r < \nu_s \left[ 1 + f(T_e/T_i) \right] \],

(14)

\[ f = \left( \frac{m_i}{m_e} \right)^{1/2} \left( \frac{T_e}{T_i} \right)^{3/2} \exp \left[ - \frac{3}{2} - \frac{T_e}{2T_i} \right] \].

(15)
Now, for zero net current \( n_r = F_{Ei} \), where the electron beam spectrum extends down to some effective cutoff energy \( E_1 \) — i.e., where the beam distribution merges with the thermal background — \( E_1 \) is a crucial parameter here, but is unfortunately not well determined by the observations. Though the ISEE power-law spectrum extends down to 5 keV, some part of this is contributed by the cooler soft-X-ray flare plasma. The later soft-X-ray maximum of the October 5 event, observed from the direction of Earth, was classified as a weak (class C) event of emission measure \( EM \approx 2 \times 10^{44} \text{ cm}^{-3} \) and temperature \( T_i = 1.6 \times 10^5 \text{ K} \) (S. R. Kane, personal communication). This (later) soft-X-ray flux \( dJ/d\epsilon \approx 4.1 \times 10^{-8} \text{ cm}^{-2} \text{ s}^{-1} \text{ eV}^{-1} \) (with \( \epsilon \) in keV) has a value of 220 cm\(^{-2}\) s\(^{-1}\) keV\(^{-1}\) at 5 keV, equal to the peak impulsive ISEE flux (243 \pm 15), and 16–37 cm\(^{-2}\) s\(^{-1}\) keV\(^{-1}\) at 7–8 keV, about half the peak impulsive ISEE flux (61.2 cm\(^{-2}\) s\(^{-1}\) keV\(^{-1}\) at 7.3–8.6 keV). These figures represent overestimates of the thermal X-ray contribution at the impulsive peak we are considering. On the other hand, a further impulsive phase thermal contribution to ISEE may arise from material hotter than the mean temperature of the gradual soft-X-ray plasma. Thus our best guess at the lower bound to the electron spectrum is \( E_1 > 7.5 \text{ keV} \). Then with \( n_r = F_{Ei} \), we need, for stability, by equations (13) and (14):

\[
A_{19}n_rT_i^{1/2} \geq 12.3/(1 + f) \quad [\text{using } F_{Ei}/F_{10} = \xi_{Ei}/\xi_{10} = (E_1/10)^{-3}] .
\]

Since \( A_{19} \approx 1 \) and \( T_i \approx 4 \) represent the extreme upper limits of plausible parameters in any flare, we see from equation (16) that for any plausible \( f \) (see below) the mean density \( n \) near the acceleration region must be significantly below \( 10^7 \text{ cm}^{-3} \) for a stable beam. In fact, we are able to test equation (16) further by utilizing the estimated soft-X-ray parameters. The small EM value probably indicates that most of the soft-X-ray source was also occulted. If we assume that the unocculted parts of the soft- and hard- (ISEE) X-ray sources occupied the same volume \( A\Delta z \), then the thermal emission measure of the hot background plasma should be (theoretically):

\[
EM = n^2A\Delta z = nA\Delta N = n_0A_{19} \times 2.5 \times 10^{46} \text{ cm}^{-3} ,
\]

where we have used equation (10). From our rough estimates of the thermal parameters \( T_i \) and EM, we have \( n_0A_{19} = 0.8 \), and, so,

\[
A_{19}n_rT_i^{1/2} \approx 1 .
\]

Comparison of the observed value (18) with the stability criterion (16) shows that (16) is satisfied if \( f \geq 11.3 \), which is close to the maximum possible value of \( f(\sim 11.12) \) at \( T_i/T_i = 3 \), suggesting that (within the considerable uncertainties in the data) the return current is not far off instability. In fact, if \( T_i/T_i \) were determined by ion-acoustic heating driven by a marginally stable return current, then, using quasi-linear resonance theory (e.g., Sagdeev and Galeev 1969), it is readily shown that \( T_i/T_i \approx 7 \) and so \( f \approx 5.4 \). The resulting factor-of-2 violation of (16) for this \( T_i/T_i \) is well within the uncertainties of our data analysis, in particular the value of \( \xi_{Ei} \) and the assumption of co-spatiality of the soft- and hard- (thin target) X-ray sources, so that the data are consistent with a marginally stable situation.

### III. Discussion and Conclusions

From § II we see that the complete set of X-ray data for the 1978 October 5 event is consistent with the spatial distributions of flux and spectrum expected from a descending electron beam model (Brown and McClymont 1975). Furthermore, the beam is accelerated at a height \( z_0 < 50,000 \text{ km} \) in an ambient medium of density \( (n \geq 10^6 \text{ cm}^{-3}) \) and temperature, which are (given the observational uncertainties), around the values necessary to satisfy the requirement of microscopic stability of the return current. This result lends support to the suggestion of Brown and Melrose (1977) that return-current stability may be the factor controlling the maximum flux in such accelerated beams, and our data analysis gives a direct constraint on conditions at the acceleration site, whereas Brown and Melrose simply used plausible values. We note that this conclusion is independent of the actual area \( A \) and ambient density \( n \) of the acceleration region since these parameters occur in the same way in the hard- and soft-X-ray data (once \( \Delta N \) has been eliminated). The injected area \( A \) may in practice be much less than \( 10^{18} \text{ cm}^2 \), in which case \( n_0 \) would be proportionately larger (e.g., \( A < 10^{18} \text{ cm}^2 \) \( \Rightarrow n < 8 \times 10^6 \text{ cm}^{-3} \) and \( z_0 \approx 28,500 \text{ km} \)). The ambient plasma electron temperature near the acceleration site is \( ~ 1.6 \times 10^7 \text{ K} \).

Although the return current is microscopically stable, to ensure that our assumption of collision domination of the beam current holds, we must also compare the collisional losses with losses under the action of the electric field driving the return current against ohmic losses (e.g., Hoyng, Knight, and Spicer 1978). Since there is no microscopically stable instability, we may take the resistivity \( \eta \) to be classical, \( \eta = \eta_0 T_i^{-3/2} \), where \( \eta_0 \approx 1.6 \times 10^{17} \text{ (Spitzer 1962)}. \) Then the ratio of ohmic to collisional losses in the thin target region (with Coulomb logarithm \( \Lambda \approx 20 \)) is

\[
\beta = \frac{\eta}{2\pi^\delta \Lambda \ln(\delta - 1)} F_i/\delta E_i = \frac{\eta_0 \xi_{Ei} E_i \delta}{2\pi^\delta \Lambda \ln(\delta - 1)} = \frac{1.0}{A_{19}n_rT_i^{3/2}} \left( \frac{E_1}{7.5} \right)^{-2} ,
\]

and from the observed values of \( T_i \) and \( A_{19}n_rT_i^{3/2} \), we find \( \beta \approx 1.0 \) for \( E_1 \approx 7.5 \text{ keV} \). Thus ohmic losses, which we have neglected thus far, may in fact be comparable to collisional losses near the acceleration site but are not so large as to violate our analysis as a first approximation within the other uncertainties involved. The effect of ohmic losses
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on the electron flux is greatest at the very lowest electron energies and is localized, rather near the acceleration site. (For a fuller discussion of these effects, see Emslie 1980.) Of course if the beam becomes unstable, or marginally so, n and thus β will increase, and anomalous ohmic losses may greatly modify the beam dynamics near the acceleration site, though not in the chromosphere.

It is also appropriate to consider the self-consistency of our model parameters in terms of energy balance. In particular we note that the beam will deposit energy collisionally in the thin target region at a rate

$$\epsilon_c = 2\pi e^2 A \frac{\delta - 1}{\delta} \left( \frac{E_{10}}{E_{10}} \right) \Delta N \frac{E_4}{E_{10}} \rho^{5/3},$$

(20)

which, for the values of $F_{10}$, $\Delta N$, and $\delta$ already used, and $E_4 = 7.5$ keV, gives $\epsilon_c = 3.7 \times 10^{27}$ ergs s$^{-1}$. This is of course much greater than the total thin target radiative loss rate ($\sim 3 \times 10^{24}$ ergs s$^{-1}$), so that if very rapid heating rates are to be avoided, thermal conduction must carry the deposited energy away to the chromosphere. The maximum (saturated) energy transport rate by conduction is

$$q_{\text{sat}} \approx 1.7 \times 10^{28} n_6 A_{19} T_7^{5/2} \text{ ergs s}^{-1},$$

(21)

or $2.7 \times 10^{28}$ ergs s$^{-1}$ for the $n_6 A_{19}$ and $T_7$ values used above. This is ample (again within the observational uncertainties) to balance the deposition, even allowing an additional return current ohmic input comparable to $\epsilon_c$ (see eq. [19]) and smaller values of $n_6 A_{19}$ and $T_7$ in the impulsive phase.

Finally we note that the consistency we have found between the descending beam model and the data here does not preclude the consistency of alternative models and in particular the dissipative thermal model (c) (Brown, Melrose, and Spicer 1979; Brown, Craig, and Karpen 1980). We consider this possibility elsewhere (Brown and Hayward 1980). In the case of the trap-plus-precipitation model (b) (Kane 1974; Melrose and Brown 1976), if the chromospheric precipitation of electrons is due to collisional pitch-angle diffusion, the ratio of coronal to chromospheric hard-X-ray emissions should be about 0.5, much larger than observed. Thus the data could only be consistent with the model if most of the arch lay below the occulting limb (25,000 km), thus hiding the bulk of the (coronal) trap component of the emission from ISEE. On the other hand, data from other flares may suggest an interpretation other than that appropriate to the October 5 event.

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REFERENCES


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