REVERSE CURRENT IN SOLAR FLARES

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ABSTRACT

We examine the proposal that impulsive x-ray bursts are produced by high-energy electrons streaming from the corona to the chromosphere. We find that the currents associated with these streams are so high that either the streams do not exist or their current is neutralized by a reverse current.

Analysis of a simple model indicates that the primary electron stream leads to the development of an electric field in the ambient corona which (a) decelerates the primary beam and (b) produces a neutralizing reverse current. It appears that, in some circumstances, this electric field could prevent the primary beam from reaching the chromosphere. In any case, the electric field acts as an energy exchange mechanism, extracting kinetic energy from the primary beam and using it to heat the ambient plasma. This heating is typically so rapid that it must be expected to have important dynamical consequences.

I. Introduction

Much of the line emission from solar flares must originate in plasma of chromospheric density. On the other hand, the corona offers a much larger volume for the storage of magnetic energy than does the chromosphere. This has led to models in which free magnetic energy, released in the corona, propagates to the chromosphere in the form of streams of high-energy particles (Sweet, 1969; Sturrock, 1974). Observational evidence for high-energy electron streams comes from analysis of microwave bursts and impulsive x-ray bursts (Kane, 1972). Evidence concerning the origin of impulsive x-ray bursts is conflicting and it is possible that some radiation originates at the chromosphere, as analyzed by Brown (1972), Hudson (1972) and Petrosian (1973), and that some originates in the corona as proposed by de Jager and Kundu (1963), Kane and Anderson (1970) and Kane (1972).

It has recently occurred to us, and to other solar physicists including Brown (1976) and Colgate (1976), that the electrical currents associated with the proposed electron streams are so high that the streams either do not exist or are neutralized by reverse currents in the ambient plasma. The reason for this statement are discussed briefly in Section II. An idealized model in which the primary current is neutralized by a reverse current is presented in Section III and applied to solar flares in Section IV. Some of the consequences of these considerations are discussed in Section V.

II. Magnetic Field Associated with Proposed Electron Stream

If energy released in the corona propagates to the chromosphere in the form of high-energy electrons, this electron stream will constitute an electrical current. If N is the total number of electrons streaming downward over the duration τ (sec.) of the impulsive phase, the magnitude J (e.m.u.) of the current may be estimated from

$$J \sim e\, c^{-1}\, N\, \tau^{-1} \quad . \tag{2.1}$$

If the transverse and longitudinal dimensions of the stream are of order L (cm.), an estimate of the strength B (gauss) of the magnetic field produced by the stream is given by

$$B \sim 2J\, L^{-1} \tag{2.2}$$

and the total energy U (erg) of this magnetic field may be estimated from

$$U \sim \frac{1}{8\pi} L^3\, B^2 \sim \frac{1}{2\pi} J^2\, L \tag{2.3}$$

which becomes

$$U \sim \frac{1}{2\pi} e^2\, c^{-2}\, N^2\, \tau^{-2}\, L \sim 10^{-40.4}\, N^2\, \tau^{-2}\, L \quad . \tag{2.4}$$

Kane and Anderson (1970) estimate the total energy involved in a typical small flare to be $\sim 10^{29}$ erg, the time scale to be $\sim 10^2$ sec, the characteristic length scale to be $\sim 10^{8.5}$ cm, and infer from the x-ray data that the total number of energetic electrons is $\sim 10^{35}$. For these values, the above formulas lead to estimates $J \sim 10^{13.2}$, $B \sim 10^5$ and $U \sim 10^{34}$. It is clear that this model, involving energy transport by

an unneutralized electron beam, leads to unacceptably high values of the magnetic field of the beam and magnetic energy associated with the beam. When one considers large flares rather than small flares, one faces the same difficulty.

One resolution of this difficulty is to abandon the assumption that energy is transported from the corona to the chromosphere by an electron stream. One might assume, for instance, that energy is transported as heat flux. The above parameters for a small flare then indicate that the mean heat flux F (erg cm^{-2} s^{-1}) would be $\sim 10^{10}$. From the formula appropriate for a fully ionized plasma (Spitzer, 1962),

$$F = \kappa \frac{dT}{dz} \sim 10^{-6} \, T^{5/2} \frac{dT}{dz} \quad , \tag{2.5}$$

we find that the electron temperature T (^{o}K) must be at least 10^{7} to give the required heat flux. For the concept of heat flux to be valid, the mean free path must be small compared with the length scale for temperature variation, e.g. 10^{-1} L, which is $\sim 10^{7.5}$ for the case now being considered. Since the mean free path γ for electron-electron and electron-ion collisions is given by

$$\gamma \approx 10^{4} \, T^{2} \, n^{-1} \quad , \tag{2.6}$$

we see that the electron density n (cm^{-3}) must be at least $10^{10.5}$. This is substantially higher than is typical in the normal corona at the initiation of a flare, although densities of this magnitude do build up in the course of a flare. This indicates that, although heat conduction may be an important, and perhaps the dominant, mechanism of energy transfer in the later stages of a flare, it cannot be the principal

mode of energy transfer at flare onset. (One should, however, bear in mind the possibility that instabilities may lead to modification of the transport properties of the plasma, modifying both the heat conduction and effective mean free path.)

III. Model of Electron Stream and Reverse Current

It is well known that a plasma tends to preserve charge neutrality. A process which tends to give an excess positive or negative charge in some region will lead to electric fields acting upon the plasma. Movement of electrons in response to this electric field will then restore charge neutrality. One expects that analogous processes will also tend to maintain current neutrality. If an electron beam is suddenly introduced into a plasma, there is a sudden change in the magnetic field structure which will develop induced electric fields opposing the primary current. For this reason, we should investigate the possibility that a reverse current develops in the corona balancing (or almost balancing) the primary current due to an assumed electron stream flowing from the corona to the chromosphere.

Our considerations are substantially simplified if we consider a symmetrical flux tube rooted in the photosphere and extending up into the corona, and assume that electrons are accelerated at the top of the flux tube by the development of stochastic electric fields or by some other mechanism. The ejection of these electrons down towards the chromosphere then leads to a charge imbalance which leads to the development of an electric field of such a magnitude that charge balance will be restored. If there is no net change in charge density, whatever current flows along the flux tube must be constant along the length of the tube. If the tube is symmetrical, this current must be zero. Hence the strong tendency of plasma to remain charge neutral implies that any primary current will generate a neutralizing secondary reverse current.

In order to simplify this exploratory analysis, we consider only a steady-state situation. We also assume that the background plasma can be adequately described by a Maxwellian velocity distribution and use transport coefficients based on this assumption (Spitzer, 1962).

We are interested in the case that the primary electron stream is composed of high-energy electrons which consequently have long mean free paths. However, we shall find that an electric field develops which decelerates the electron stream. But when the electron energy becomes comparable with the thermal energy, the mean free path will be sufficiently short that the primary electrons will merge with the background plasma. As a simple representation of this process, we ignore collisions in discussing the primary beam but we assume that an electron of the primary beam is absorbed into the background plasma when it is decelerated to zero energy.

If, as a further simplification, we consider a magnetic flux tube of constant cross section, we may use the following simple one-dimensional form of the Vlasov equation:

$$v\frac{\partial f}{\partial s} + \frac{e}{m}\frac{d\varphi}{ds}\frac{\partial f}{\partial v} = 0 \qquad (3.1)$$

where s measures arc length along the tube, v is velocity (along the tube), $f(s,v)$ is the velocity distribution function of the primary electron stream, and φ is the electrostatic potential.

At the top of the loop ($s = 0$), the primary electron stream is moving with positive velocity on the side corresponding to increasing s. Hence we may, without ambiguity, express f in terms of ψ defined by

$$\psi = \frac{mv^2}{2e} \quad . \qquad (3.2)$$

The initial distribution function may therefore be expressed as

$$f(0, v) = F(\psi) \quad . \tag{3.3}$$

With this initial condition, we find that the solution of the Vlasov equation (3.1) is

$$f(s, v) = F(\psi - \varphi) \quad . \tag{3.4}$$

The current density j_s (e.m.u.) in the primary electron stream is given by

$$j_s = -\frac{e}{c}\int_0^\infty f(s, v)\, v dv \tag{3.5}$$

which may be expressed as

$$j_s = -\frac{e^2}{mc}\int_0^\infty F(\psi - \varphi)\, d\psi \quad . \tag{3.6}$$

Since φ will prove to be negative in the region of interest, it is convenient to write

$$\theta = -\varphi \tag{3.7}$$

so that (3.6) may be reexpressed as

$$j_s = -\frac{e^2}{mc}\int_\theta^\infty F(X)\, dX \quad . \tag{3.8}$$

Charged neutrality requires that

$$j_p + j_s = 0 \tag{3.9}$$

where j_p is the secondary current induced in the background plasma. We here assume that the density and temperature are such that this may be

represented by Ohm's law,

$$j_p = \sigma E = \sigma \frac{d\theta}{ds} \quad , \tag{3.10}$$

where σ is the electrical conductivity expressed in modified gaussian units.

It is convenient to introduce a new independent vairable ξ to replace s, by the relationship

$$d\xi = \frac{ds}{\sigma(s)} \quad . \tag{3.11}$$

Then, on substituting equations (3.8) and (3.10) into (3.9) and differentiating with respect to ξ, we obtain

$$\frac{d^2\theta}{d\xi^2} + \frac{e^2}{mc} F(\theta) \frac{d\theta}{d\xi} = 0 \quad . \tag{3.12}$$

It is convenient to solve this equation for ξ in terms of θ,

$$\xi = X(\theta) \quad , \tag{3.13}$$

rather than vice versa. Equation (3.13) becomes

$$\frac{mc}{e^2} \left(\frac{dX}{d\theta}\right)^{-2} \frac{d^2X}{d\theta^2} = F(\theta) \quad , \tag{3.14}$$

which may be integrated once to give

$$\frac{mc}{e^2} \left[\left(\frac{d\theta}{d\xi}\right)_{\xi=0} - \left(\frac{dX}{d\theta}\right)^{-1} \right] = \int_0^\theta F(\theta')\cdot d\theta' \quad , \tag{3.15}$$

if we assume that $\theta = 0$ ($\varphi = 0$) and $X = 0$ ($\xi = 0$) at $s = 0$. We find, from (3.8), (3.9) and (3.10), that

$$\frac{e^2}{mc}\left(\frac{d\theta}{d\xi}\right)_{\xi=0} = \int_0^\infty F(\acute{\theta})\, d\acute{\theta}^{-1} \quad . \tag{3.16}$$

Hence equation (3.15) becomes

$$\frac{dX}{d\theta} = \frac{mc}{e^2}\left[\int_0^\infty F(\acute{\theta})\, d\acute{\theta}\right]^{-1} \quad . \tag{3.17}$$

IV. Specific Model

It is now convenient to introduce a specific form for $F(\psi)$,

$$F(\psi) = K(\psi_0 + \psi)^{-\gamma} \quad . \tag{4.1}$$

This is a power-law distribution at high energy which flattens at low energy, the "knee" being characterized by ψ_0.

We introduce the symbol $H(\psi, s)$ for the flux of electrons ($cm^{-2}\ s^{-1}$) of energy exceeding $e\psi$ at the position s

$$H(\psi, s) = \frac{e}{m} \int_{\psi}^{\infty} F\left(\psi' + \theta(s)\right) d\psi' \quad . \tag{4.2}$$

If the initial flux is written as $H_0(\psi)$, we find that

$$H_0(\psi) = \frac{eK}{(\gamma-1)m} \left(\psi_0 + \psi\right)^{-\gamma+1} \quad , \tag{4.3}$$

so that the total particle flux is given by

$$H_0(0) = \frac{eK}{(\gamma-1)m} \psi_0^{-\gamma+1} \quad . \tag{4.4}$$

With the form (4.1) for $F(\psi)$, equation (3.17) integrates to give

$$X(\theta) = \frac{\gamma-1}{\gamma} \frac{mc}{e^2K} \left[\left(\psi_0 + \theta\right)^{\gamma} - \psi_0^{\gamma}\right] \quad . \tag{4.5}$$

We easily obtain from (4.5) an expression for the (negative) electric potential θ in terms of the modified distance measure ξ:

$$\theta(\xi) = \left[\psi_0^{\gamma} + \frac{\gamma}{\gamma-1} \frac{e^2K}{mc} \xi\right]^{1/\gamma} - \psi_0 \quad . \tag{4.6}$$

Hence, from (4.2), we find that

$$H(\psi, s) = \frac{eK}{(\gamma-1)m} \left[\psi + \left\{\psi_0^{\gamma} + \frac{\gamma}{\gamma-1} \frac{e^2K}{mc} \xi\right\}^{1/\gamma}\right]^{-\frac{(\gamma-1)}{\gamma}} \quad . \tag{4.7}$$

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On noting that the electric current carried by the stream is related to $H(\psi, s)$ by

$$j_s(s) = -\frac{e}{c} H(0, s) \quad , \tag{4.8}$$

we see that

$$j_s(s) = -\frac{e^2 K}{(\gamma-1)mc}\left[\psi_0^{\gamma} + \frac{\gamma}{\gamma-1}\frac{e^2 K}{mc}\xi\right]^{-\frac{\gamma-1}{\gamma}} \quad . \tag{4.9}$$

In order to specify the current, particle flux and electric field as functions of s, we must adopt a specific form for $\sigma(s)$. A convenient approximation to the structure of the solar atmosphere, which is expressible in analytic form, is provided by the constant heat flux model. If we now assume that s measures distance vertically downwards from the corona, and that $n = n_0$, $T = T_0$, at $s = 0$, this model (Adams and Sturrock, 1975) yields the following expressions:

$$T(s) = \left(T_0^{7/2} - bFs\right)^{2/7} \tag{4.10}$$

$$n(s) = n_0\left(T_0/T(s)\right) \exp\left\{-\frac{a}{F}\left[\left(T_0^{7/2} - bFS\right)^{5/7} - T_0^{5/2}\right]\right\} \quad , \tag{4.11}$$

where $a \approx 10^{-10.21}$, $b \approx 10^{6.58}$, and F (erg cm^{-2} s^{-1}) is the downward heat flux.

The electrical conductivity, in modified gaussian units, may be derived from the expression given by Spitzer (1962):

$$\sigma = g\ T^{3/2} \tag{4.12}$$

where $g \approx 10^{-3.64}$. Hence we find, from (3.11), that ξ is related to

s by

$$\xi = \frac{7}{4}\frac{1}{bgF}\left[T_0^2 - \left(T_0^{7/2} - bFs\right)^{4/7}\right] \quad . \tag{4.13}$$

Our model is then completely specified by the choice of the coronal temperature, the coronal density, the coronal heat flux, γ, the energy corresponding to ψ_0, and the injected energetic electron flux. For the coronal parameters, we have adopted values typical of the corona above an active region (Noyes, 1971):

$$T \simeq 3 \times 10^6 \text{ K}$$
$$n \simeq 10^9 \text{ cm}^{-3}$$
$$F \simeq 5 \times 10^6 \text{ erg cm}^{-2} \text{ sec}^{-1} \quad .$$

We choose ψ_0 to correspond to 25 kev and we choose $\gamma = 2.5$, corresponding to an x-ray spectral index of $\sim 3 - 3.5$ (Brown, 1975), which is a reasonable value for the spectral index at peak emission (Datlowe, 1975).

The fraction of the beam energy deposited, and the total energy deposited by joule heating between $T = 3 \times 10^6$ K and $T = 3 \times 10^4$ K as a function of the energetic electron flux, are displayed in Figure 1. For a flare area of $10^{19.5}$ cm^2, the energetic electron flux inferred from a large impulsive x-ray burst corresponds to $\sim 10^{17}$ cm^{-2} sec^{-1} (Hoyng et al., 1975). Figure 2 illustrates the energy deposition rate due to joule heating as a function of temperature of the atmosphere for this injected energetic electron flux. The ordinate of Figure 2 is the time required to raise the ambient plasma temperature by 10^7 K.

V. Discussion

For the adopted model, Figure 2 shows that the steady-state assumption is not justified for the energetic electron flux required to produce the observed x-ray flux by thick target bremsstrahlung. The atmosphere, especially the dense low-temperature region, will be heated rapidly to a high temperature. Once the plasma is raised to a few times 10^7 K, ohmic losses will become unimportant since the joule heating rate is inversely proportional to the electrical conductivity. Joule heating is a more effective mechanism for heating the plasma than collisional losses from the energetic electrons, because the ohmic losses are caused by thermal electrons in the reverse current which have much shorter mean free paths than the energetic electrons. It is likely that a larger fraction of the beam energy is deposited than is indicated in Figure 1 for two reasons. First, the constant heat flux model underestimates the extent of the low temperature region where the ohmic losses are largest. Second, if the drift velocity associated with the reverse current becomes too large, instabilities can develop and reduce the electrical conductivity of the plasma. If the conductivity of the plasma is low enough, the electric field that develops may be large enough to effectively prevent the stream from reaching the chromosphere.

The time scale for heating the dense low-temperature plasma is comparable with the time scales characteristic of impulsive x-ray bursts (Hoyng et al., 1975). It is possible that thermal bremsstrahlung from the rapidly heated plasma can account for a significant portion of the observed impulsive x-ray flux. Hence this mechanism offers an explanation of the fact that some flares first produce high-energy x-ray emission

near the top of a loop rather than at the foot-points of the loop (Brueckner, 1976). Another important consequence of this process is that, if thermal emission can account for a substantial fraction of the impulsive flux up to ~ 50 kev, then the number of electrons required to produce the non-thermal x-ray flux is greatly reduced (Brown, 1975).

The rapid heating of the plasma will also cause a large pressure imbalance. The time τ (sec.) for the plasma to respond to pressure changes by bulk motions can be estimated from

$$\tau \sim \frac{L}{V_s} \quad , \tag{5.1}$$

where L is a characteristic length and V_s is the ion thermal velocity. Even for a temperature as large as 10^7 K, this time is long (10^2 sec) compared to the heating time for a length scale of $10^{9.7}$ cm, so that the plasma density will not change appreciably during the heating. In the constant heat flux model we have adopted, the density at 3×10^4 K is approximately 100 times the density at 3×10^6 K. If both these regions are heated to a temperature of a few times 10^7 K, the resulting pressure imbalance must be expected to produce explosive evaporation of high density plasma into the corona.

A more realistic time-dependent model is necessary to account for the complicated phenomena observed in solar flares. It is also possible that collective effects are important and that the assumption of a Maxwellian velocity distribution for the ambient plasma must be abandoned. However, our analysis of reverse current in terms of an idealized model strongly suggests that joule heating is likely to play an important role in the heating of solar plasma during a flare.

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Figure Captions

Figure 1. The fraction of the beam energy deposited (solid curve) and the total energy deposited (broken curve) by joule heating as a function of the energetic particle number flux.

Figure 2. Time τ_H required for ohmic losses to heat the plasma by 10^7 K as a function of the ambient plasma temperature.

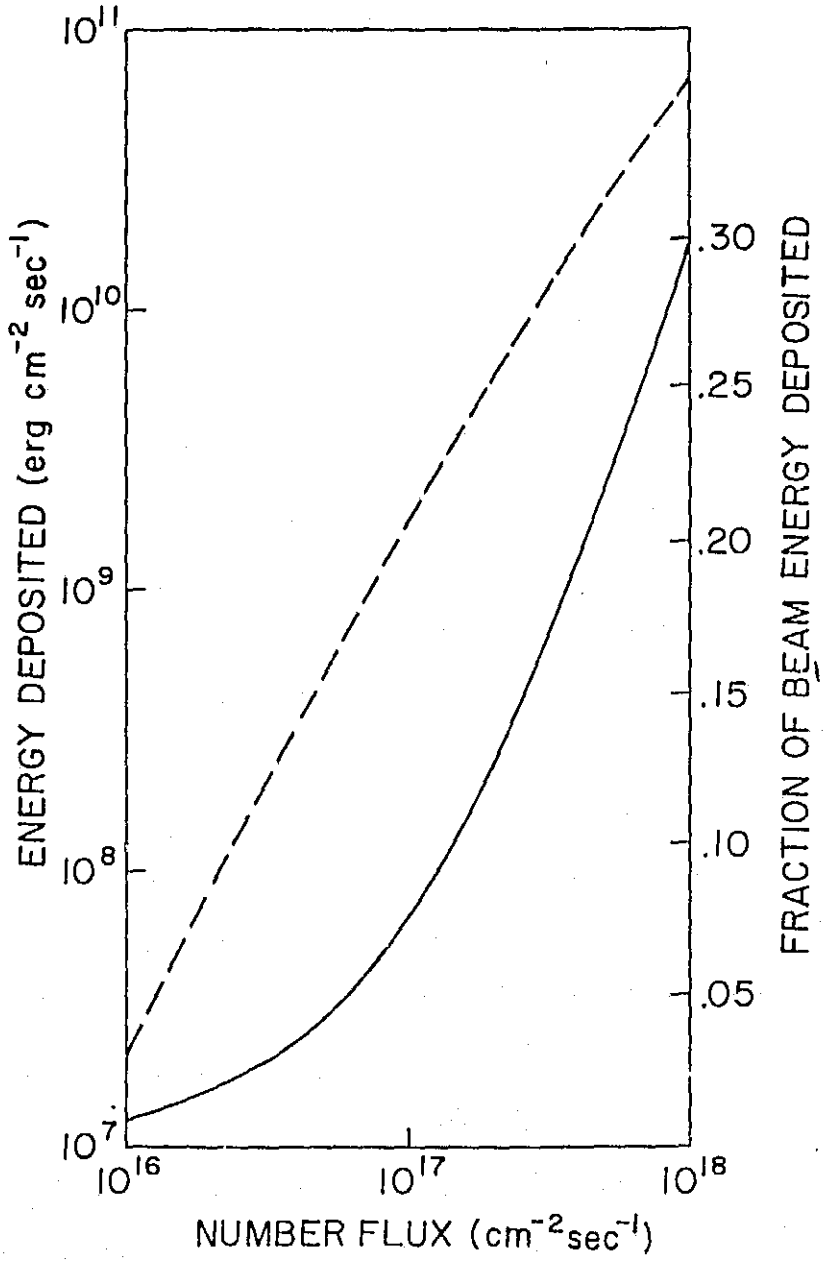
10^{11}
10^{10}
10^{9}
10^{8}
10^{7}
ENERGY DEPOSITED (erg cm^{-2} sec^{-1})
.30
.25
.20
.15
.10
.05
FRACTION OF BEAM ENERGY DEPOSITED
10^{16}
10^{17}
10^{18}
NUMBER FLUX (cm^{-2}sec^{-1})

Figure 1

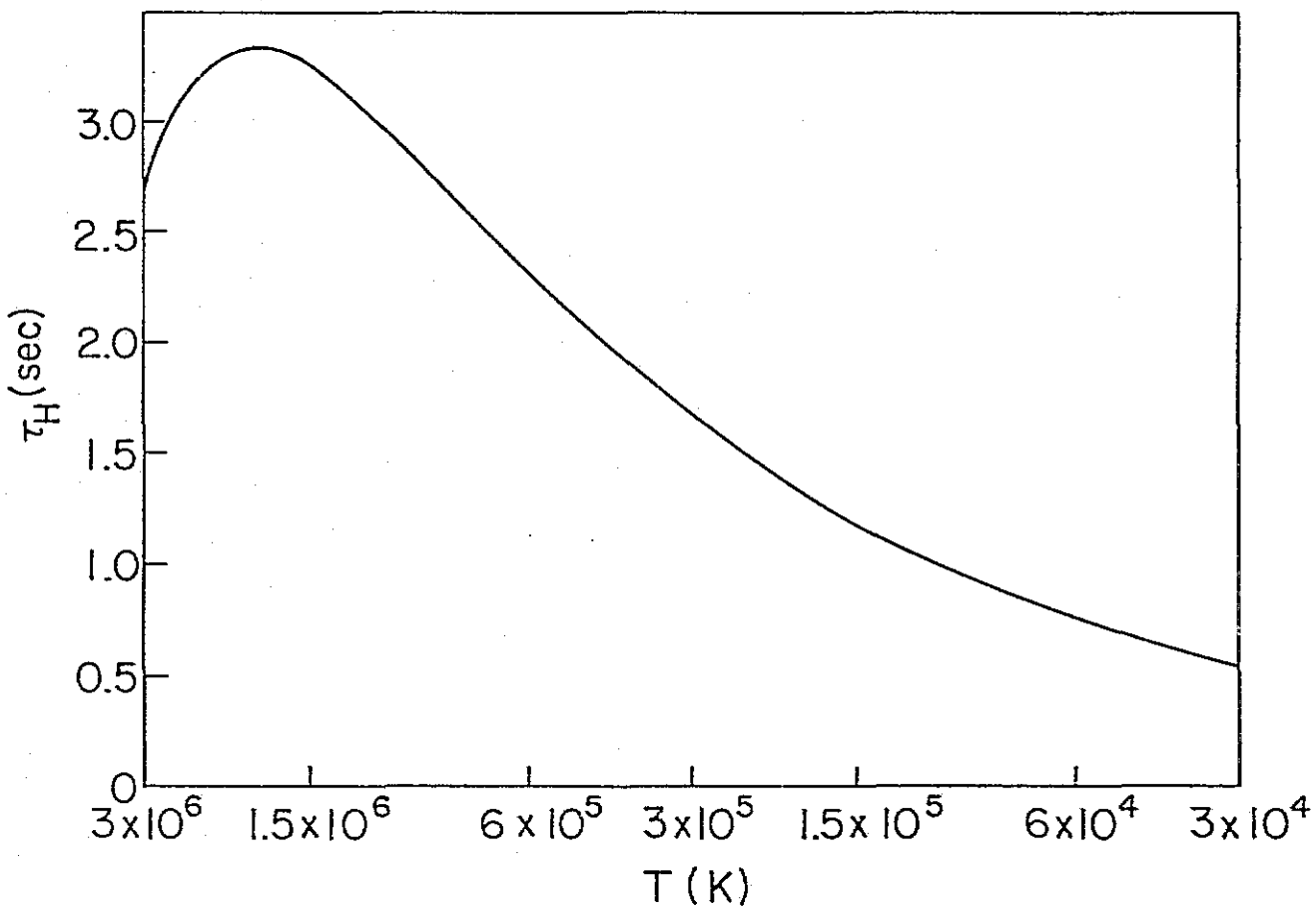
3.0
2.5
2.0
1.5
1.0
0.5
0
τ_H (sec)
3×10^6
1.5×10^6
6×10^5
3×10^5
1.5×10^5
6×10^4
3×10^4
T (K)

Figure 2