COHERENT RADIATION FROM ELECTROSTATIC DOUBLE LAYERS[1]

Jan Kuijpers
Sterrekundig Instituut Utrecht
P.O.Box 80 000
3508 TA Utrecht
The Netherlands

Abstract

The radiation properties of electrostatic double layers (DLs) are of potential significance for cosmic magnetic explosions. If the emission is sufficiently intense it can serve as a diagnostic tool for particle acceleration by localized strong potential drops in current-carrying plasmas. Moreover such intense emission may form the explanation of some of the observed intense and narrow-band bursts of radiation from stellar and planetary magnetospheres. Here we study the efficiency of two coherent radiation processes: antenna radiation and a maser process. It is found that both processes can operate in the DL to produce intense and narrow-band emission. Antenna radiation occurs if the dimensions of the double layer are smaller than the wavelength of the emitted radiation. This process is therefore relevant to laboratory rather than to astrophysical plasmas. The maser on the other hand requires an amplification length inside the double layer much larger than the emitted wavelength, and can lead to observable emission in astrophysical circumstances. The growth is exponential and the rate depends only on the electric field energy density of the DL. Since the latter is externally controlled by the electric circuit it is a constant for the emission process so as to constitute a true maser. The maximum brightness temperature is of order $10^{25} K$. Masing radiation from electrostatic DLs is therefore a candidate for some of the observed intense narrow-band cosmic radio emission.

1 Introduction

Electrostatic double layers (DLs) have first been proposed in astrophysical context by Alfvén (1958). Since then they have been advocated in a variety of cosmic high-energy phenomena (Alfvén, 1981). Recent reviews on the subject are given by Raadu and Rasmussen (1988) and Raadu (1989). DLs are of particular interest since the energy of a global electric circuit can be released locally in DLs (Jacobsen and Carlqvist, 1964) which have large voltage drops. DLs are therefore of importance for the explosive release of stored magnetic energy as is observed in *solar and stellar flares* (Kuijpers, 1989a). The important dynamic properties of DLs are the *acceleration of both electrons and ions and the triggering of mass motions*, effects which are indeed characteristic for many cosmic explosive phenomena. The particle acceleration is caused by the electric potential drop parallel to the magnetic field; the mass motion in the ambient ideal (that is non-resistive) mhd plasma is associated with a perpendicular potential difference, which again is connected to the parallel potential drop in a realistic finite 3-D structure (Raadu, 1984).

On the theoretical side much effort has been put into the construction of stationary DLs. A stationary DL can be considered as a coherent large-amplitude plasma wave: The electric field derives from a charge separation which is maintained self-consistently by the particle motion under the influence of the electric force. Laboratory experiments show that in practice stationary DLs are rather difficult to maintain and that "*explosive*" DLs are the rule.

Of particular interest for astrophysics are the *strong DLs*, in which the energy of an accelerated particle is much larger than the thermal energy per particle outside the DL. If indeed strong DLs are the locations of energy release in cosmic magnetic explosions the question arises how they can

[1]The calculations of the results presented in this paper can be found in Kuijpers 1989b.

E. R. Priest and V. Krishan (eds.), Basic Plasma Processes on the Sun, 365–373.

be observed. In this paper we calculate the coherent emission from DLs. As a first approach we shall restrict ourselves to ***stationary and strong*** DLs. For our model calculations we use the simple configuration of Fig. 1 and write the potential jump ϕ as

$$e\phi = uKT_0, \tag{1}$$

with $u \simeq 10-10^4$, T_0 the temperature outside the DL, e the electron charge and K is Boltzmann's constant. For the characteristic thickness of the DL along the electric field direction, taken to coincide with the magnetic field direction, we write

$$L = \ell\lambda_{Do}, \tag{2}$$

with $\ell \simeq 10-10^3$ and λ_{Do}, the Debye wavelength outside the DL. The shape of the DL is modelled as a cylinder with radius R. By assumption the electric field is confined to the DL, it is uniform within the DL and directed parallel to the cylinder axis.

In Section 2 we calculate the coherent emission from a ***non-relativistic*** DL. In Section 3 we apply the results of the calculation to laboratory experiments and cosmic plasmas. The conclusions are given in Section 4.

2 Coherent Emission

As a charged particle inside a DL is accelerated it emits electromagnetic radiation. For a single particle the power radiated is

$$P = \frac{2q^2\gamma^2}{3m^2c^3}\left\{\left|\frac{\mathrm{d}\vec{p}}{\mathrm{d}t}\right|^2 - \frac{1}{c^2}\left(\frac{\mathrm{d}U}{\mathrm{d}t}\right)^2\right\}, \tag{3}$$

with q the charge, m the mass, γ the Lorentz factor, $\vec{p}$ the momentum and U the energy of the particle. Inside the DL the force on the particle and its rate of change of energy are given by ($\vec{E}$ is the electric field)

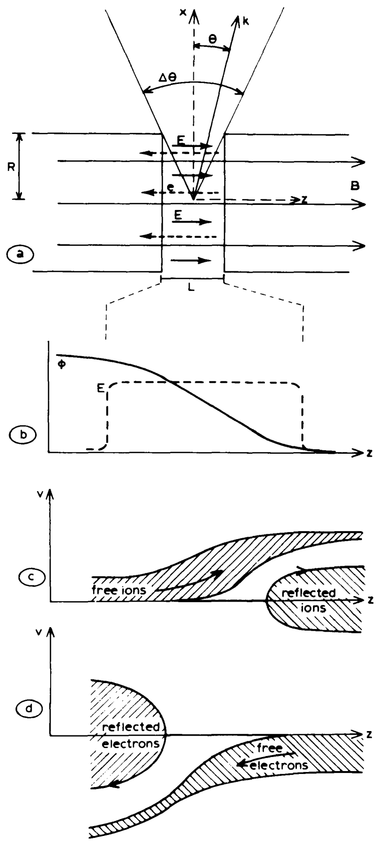
$$\begin{aligned}\frac{\mathrm{d}\vec{p}}{\mathrm{d}t} &\simeq q\vec{E},\\ \frac{\mathrm{d}U}{\mathrm{d}t} &\simeq \frac{\mathrm{d}p}{\mathrm{d}t}v,\end{aligned} \tag{4}$$

since for a strong DL the particle speed v is largely parallel to the electric and magnetic fields. With these expressions Eq. (3) reduces to

$$P = \frac{2q^2}{3m^2c^3}\left|\frac{\mathrm{d}\vec{p}}{\mathrm{d}t}\right|^2 = \frac{2q^4E^2}{3m^2c^3}, \tag{5}$$

independent of the particle energy γ. Clearly this "*linear accelerator emission*" (a term introduced by Melrose (1980) in the context of radio pulsars) is a relatively inefficient process when the (single) particle is moving at relativistic speed. For comparison: the power emitted in gyrosynchrotron radiation is given by a relation similar to Eq. (5) with E replaced by $\gamma\beta B_\perp$ ($\beta \equiv v/c$ and $B_\perp$ is the magnetic field component perpendicular to the particle velocity v). Further Eq. (5) shows that electrons radiate much more efficiently in a DL than ions, as often happens in radiation processes. In principle the efficiency can be greatly enhanced if the particles radiate in phase *(Antenna radiation)* or if the radiation is amplified *(Maser)*. We shall investigate both possibilities in turn for non-relativistic DLs.

2.1 Antenna radiation



When the wavelength of the emitted radiation is larger than the dimensions of the emitting ensemble of electrons, an enhancement of the emission occurs in comparison with the case of an ensemble of independent radiators: in the former case the amplitudes of the emitted waves add coherently. The increase of the emission with respect to the incoherent case is expected to be a factor N, the number of particles in the bunch. Indeed such a result obtains from a rigorous calculation of the emission (Kuijpers, 1989b).

Let us consider the radiation from an ensemble of electrons (charge e and mass m). As the electric field is assumed to be confined to the DL (see Fig. 1) also the radiation originates from within the DL.

Figure 1: (a) shows the geometry of the DL with dimensions L and R parallel and respectively perpendicular to the magnetic field $\vec{B}$, the electric field $\vec{E}//\vec{B}$, the electron beam, and the wave vector $\vec{k}$ of escaping radiation. The lower figures show from top to bottom: the electric potential and field in the DL (b), level curves of the ion distribution function (c) and of the electron distribution function (d) with v the velocity in the z-direction.

For a strong DL the electron transit time t is practically independent of the initial electron speed:

$$t \simeq \left(\frac{2mL}{eE}\right)^{1/2} = \frac{\ell}{\omega_{po}}\left(\frac{2}{u}\right)^{1/2} \equiv \frac{\epsilon^{1/2}}{\omega_{po}}. \tag{6}$$

ω_{po} is the electron plasmafrequency outside the DL, m is the electron mass and the last equality defines ϵ. We choose the $\hat{z}$-direction to coincide with the direction of the background magnetic and electric field, and the ($\hat{x}$,$\hat{z}$)- plane to contain the wave vector $\vec{k}$. Since the acceleration occurs in the $\hat{z}$-direction, only the component j_z contributes to the emission. Then only the wave mode with polarization vector in the $(\vec{k},\vec{B})$-plane, the so- called ordinary mode, is radiated. Further we use the approximations

$$\begin{aligned} 2kR\cos\theta &\leq 1, \\ kL\sin\theta &\leq 1, \end{aligned} \tag{7}$$

where θ is the angle between the $\vec{k}$ and $\hat{x}$- direction. Condition (7b) guarantees that the contributions to the Fourier transform of the current density arising from one electron at the different z-positions along its orbit, add constructively, while condition (7a) ensures that the contributions from different x-coordinates add constructively. Since the DL is assumed to be strong, the transit time is independent of the initial particle speed, and therefore also of the initial spread in velocities. We approximate $\omega^t(\vec{k}) = (\omega_{po}^2 + k^2c^2)^{1/2}$ for $\omega_{co} << \omega_{po}(\omega_{co}$ is the electron cyclotron frequency,

ω_{po} the electron plasmafrequency, both taken outside the DL). Then the energy radiated per unit frequency, solid angle and time $p^t(\omega)$ (the differential power) for a stationary DL is

$$p^t(\omega) = \frac{N^2 e^4 E^2 \epsilon^{1/2} \cos^2\theta}{\pi^2 m^2 c^3 \omega_{po}} \frac{(1-\cos x)^2}{x^4} \left(1 - \frac{\epsilon}{x^2}\right)^{1/2}. \tag{8}$$

Here we use the abbreviations $x \equiv \omega t, \epsilon \equiv 2\ell^2/u = (\omega_{po} t)^2$ (see Eq. (1) and Eq. (2)).

2.1.1 Frequency, Bandwidth and Polarization

To determine the peak and bandwidth of the radiated power we approximate the instantaneous number of electrons in the DL with

$$N = A n_o v_{to} t, \tag{9}$$

where A is the cross-section of the DL perpendicular to the magnetic field n_o is the density of the ambient plasma and $v_{to} \equiv (KT_o/m)^{1/2}$, and we use $E \simeq \phi/L$ and Eq. (1) and Eq. (2), so that Eq. (8) can be rewritten as

$$\begin{aligned} p^t(\omega)\omega_{po} &= \frac{A^2}{(\ell\lambda_{Do})^4} \frac{(KT_o)^4 \cos^2\theta}{2^5 \pi^4 m^2 c^3 e^2} u^3 G(x), \\ \text{with the definition } G(x) &\equiv \frac{\epsilon^{5/2}(1-\cos x)^2}{x^4} \left(1 - \frac{\epsilon}{x^2}\right)^{1/2}. \end{aligned} \tag{10}$$

The function $u^3 G(x)$ is a measure of the power radiated by one DL. For different values of ϵ Table 1 gives the relative peak frequency $\omega/\omega_{po} = x\epsilon^{-1/2}$ (see Eq. (6)), the corresponding primary maximum of $G(x)$ and the relative half width $\Delta\omega/\omega_{po}$ of the primary maximum. Typical values for ϵ are in the range $2 - 2\ 10^4$ as u varies from $10 - 10^4$ and ℓ from $10 - 10^3$.

TABLE 1

Peak and halfwidth of the dimensionless power function $G(x)$
$x \equiv \epsilon^{1/2}\omega/\omega_{po}$

ϵ	2000	200	20	2	0.2	0.02
ω/ω_{po}	1.19(1)	1.11	1.04(2)	1.34	2.09	3.57
$\Delta\omega/\omega_{po}$	0.043	0.13	0.12	0.80	1.92	3.84
max$G(x)$	48	16	1.2	0.51	$3.4\ 10^{-3}$	$1.3\ 10^{-5}$

(1) Secondary maxima above half the primary maximum exist both at lower and higher frequencies.
(2) A secondary maximum above half the primary maximum exists at a higher frequency.

From the table and Eq. (10) it follows that for a fixed temperature of the ambient plasma and a fixed shape of the DL the maximum radiation is emitted for strong (large u) and extended (large ℓ) DLs above the plasmafrequency of the ambient plasma ω_{po} and has a characteristic bandwidth of 10 % .

The emission at the source is completely linearly polarized and in the ordinary mode (electric polarization in the plane $(\vec{k}, \vec{B})$).

2.1.2 Brightness Temperature

The radiated power must always be smaller than the rate at which the electric current energy is transformed into particle acceleration:

$$\int\int p^t(\omega)\mathrm{d}\omega\mathrm{d}\Omega < I\phi \simeq An_o v_{to} e\phi = An_o v_{to} uKT_o$$

or using Eq. (10),

$$\frac{2}{3\pi^2}\frac{A}{(\ell\lambda_{Do})^2}\left(\frac{v_{to}}{c}\right)^3\frac{u}{\epsilon^{3/2}}\int_o^\infty G(x)\mathrm{d}x < 1. \tag{11}$$

Condition (11) puts an upper limit to the lateral cross section A of the DL.

An estimate of the brightness temperature T_b of the power calculated in Eq. (10) is obtained from the definition of $T_b(f)$ at a frequency $f = \omega/2\pi$:

$$\frac{f^2}{c^2}KT_b(f) \equiv \frac{p^t(\omega)\mathrm{d}\omega}{A(\theta)\mathrm{d}f}, \tag{12}$$

where $A(\theta)$ is the projected area of the DL perpendicular to the direction θ. Substituting Eq. (10) into Eq. (12) one finds

$$T_b(f) = T_o N_{Do}\frac{u^3}{\ell^2}\frac{v_{to}}{c}\frac{A^2}{(\ell\lambda_{Do})^4}\frac{(\ell\lambda_{Do})^2\cos^2\theta}{A(\theta)}\frac{G(x)}{(\omega/\omega_{po})^2} \tag{13}$$

with $N_{Do} \equiv n_o\lambda_{Do}^3$ the Debye number of the ambient plasma. At the intensity maximum ($2\pi f = \omega \simeq \omega_{po}$ and $x \simeq \epsilon^{1/2}$) the brightness temperature is enhanced with respect to the ambient temperature by a factor $N_{Do}u^3\ell^{-2}v_{to}c^{-1}$ (e.g with a factor 4 10^8 for $T = 10^5 K, n = 10^{10} cm^{-3}, u = 10^4, \ell = 10^3$).

Finally the brightness temperature (13) can never exceed the optically thick limit, given by the "effective temperature" of the coherent bunch of particles:

$$T_b(f) \leq NuT_o = T_o N_{Do} u^{1/2}\ell^3\frac{A}{(\ell\lambda_{Do})^2} \tag{14}$$

2.2 Maser

In the previous section we have considered the possibility of coherent emission from a DL in the form of antenna radiation. A necessary condition for the antenna process is that the DL should be small in comparison to the wavelength of the emission. In this section we consider the possibility of amplification of radiation in the DL in the opposite case when the lateral dimensions (perpendicular to $\vec{E}$ and $\vec{B}$) of the DL are large in comparison with the wavelength. A priori it is clear that the extension of the DL along $\vec{E}$ is always smaller than the wavelength of the radiation, since

$$k^t < \omega^t/c = O(2\pi/t)/c \simeq O\left(\frac{\pi v_f}{Lc}\right) < \frac{\pi}{L}, \tag{15}$$

where we have used Eq. (6) for the transit time of an electron in the DL and v_f is the fast electron speed obtained in a strong DL.

The process which we are looking for produces electromagnetic waves from the interaction between the electron beam created in the DL and the electric field inside the DL. Now *the electric field inside a strong DL is strong in several respects*: Firstly in a strong DL it is strong with respect to the critical runaway field ($E_c = 0.16E_D, E_D$ is the Dreicer field) in the ambient plasma,

$$\frac{E_{DL}}{E_c} \simeq \frac{\phi\lambda_{Do}^2}{Le0.16\ell n\Lambda} = \frac{u}{\ell}\frac{4\pi N_{Do}}{0.16\ell n\Lambda} >> 1. \tag{16}$$

https://doi.org/10.1017/S0074180900088276 Published online by Cambridge University Press

Secondly the energy density in the electric field inside a strong DL is at least comparable to the thermal energy density of the ambient plasma

$$\frac{E^2/8\pi}{3n_oKT_o} = 4\pi\left(\frac{u}{\ell}\right)^2. \tag{17}$$

Here we shall consider only weakly nonlinear effects as a first approach; It is expected however that the efficiency will be comparable or higher in a realistic DL. In this approximation the simplest interactions which produce electromagnetic waves, are *1. the scattering of the electric field of the DL off the fast electrons into an escaping high-frequency electromagnetic wave, and 2. the related process of double emission of DL field and electromagnetic field.* The process works if a particular resonance condition is fulfilled: the Doppler shifted (in the frame of the fast electron) frequency must equal plus (for scattering) or minus (for double emission) the Doppler shifted frequency of the electromagnetic wave

$$\omega^{DL}\left(\vec{k}^{DL}\right) - \vec{k}^{DL}\cdot\vec{v} = \pm\left(\omega^t(\vec{k}^t) - \vec{k}^t\cdot\vec{v}\right). \tag{18}$$

The electric field of the DL is static or slowly-changing on an electron transit time so that we can approximate $\omega^{DL} = 0$. Further, substantial amplification of radiation can only take place transverse to the applied electric field (cf. Eq.(15)). Eq.(18) then reduces for a nonrelativistic DL to

$$\omega^t \simeq k^{DL}v_f \simeq 2^{3/2}\pi u^{1/2}\omega_{po}/\ell, \tag{19}$$

where v_f is the electron beam velocity at the end of the DL and we have used $k^{DL} \simeq 2\pi/L$. From Eq.(19) it is clear that scattered radiation above the ambient plasmafrequency is possible. In the scattering process the small frequency of the DL electric field is upconverted by the energetic electrons into the relatively high frequency of the radiation. In a thermal plasma scattering of low-frequency waves off electrons is usually an inefficient process since the induced electron (or Thomson) scattering is practically cancelled by the simultaneous nonlinear scattering on the polarization clouds (opposite sign) of the electrons. Here however the electrons acquire a beamlike distribution in the strong DL, they loose the company of polarization clouds and induced Thomson scattering becomes important.

We have calculated the efficiency of the scattering (and double emission) process in a DL using the *weak turbulence approximation*. The calculations are presented elsewhere (Kuijpers, 1989b). Using an electron beam distribution function

$$f_o(\vec{p}) = \frac{n_f}{m}(2\pi)^{-3/2}v_{to}^{-3}\exp\left\{-[(v_z+v_f)^2+v_x^2+v_y^2]/2v_{to}^2\right\} \tag{20}$$

the calculated *growth rate* is

$$\gamma^t(\vec{k}^t) = \frac{u^{3/2}}{\ell^2}\frac{\omega_{po}^4}{\omega_t^3}y exp(-y^2/2), \tag{21}$$

where $y \equiv (v_f - \omega_t/\mid k^{DL}\mid)/v_{to}$ and we have used $\frac{1}{2}mv_f^2 = e\phi, W^{DL} \simeq \phi^2/8\pi L^2, n_fv_f = n_ov_{to}$. It is important to realize that this form of coherent radiation from a DL constitutes a *true maser* since the energy density of the electric field in the DL is fixed and maintained in principle by the much larger energy of the global electric circuit of which the DL forms part.

2.2.1 Frequency, Bandwidth and Polarization

The characteristic frequency of the radiation emitted in the maser is determined by the resonance condition (18). For a non- relativistic DL the frequency is given by Eq. (19) and the ratio of the frequency with respect to the plasmafrequency increases with the strength of the DL.

The bandwidth is determined by 1. the frequency dependence of the growth rate Eq. (21): $\gamma \alpha \omega_t^{-3} y exp(-y^2/2)$, which sets an upper limit to the bandwidth of order $u^{-1/2}$ and 2. the amplification length Δx in the source:

$$\frac{\Delta\omega}{\omega} \simeq \frac{c}{3\Delta x \gamma^t} \tag{22}$$

For an e-folding length of 50 (see below) the relative bandwidth of the emission is 4 10^{-3}, much narrower than the bunched emission of Section 2.1.

The strongest radiation is emitted in directions perpendicular to the DL axis and is *completely polarized in the ordinary mode.* This follows directly from the vanishing nonlinear cubic current for the case of a DL electric field parallel to the z-axis, escaping radiation in the x-direction and electric field of the radiation parallel to the y-direction (the extraordinary mode polarization).

2.2.2 Brightness Temperature, Flux and Angular Extent

The brightness temperature of the maser depends on the number of e-folding lengths over which the radiation can be amplified inside the DL and on the absorption outside the DL. The observed flux from a stationary DL of thickness $\ell\lambda_{Do}$ and radius R at distance D from the observer is related to the brightness temperature of the radiation by

$$S = \frac{f^2}{c^2} K T_b \frac{\ell\lambda_{Do} R}{D^2} \simeq 3.5\ 10^{-34} n_9^{1/2} T_6^{1/2} \frac{\ell}{100} \frac{R}{3km} \frac{T_b}{10^6} erg\ cm^{-2} s^{-1} Hz^{-1}, \tag{23}$$

where we have neglected intervening absorption. For a flux of $10^{-17} erg\ cm^{-2} s^{-1} Hz^{-1} = 10^6 Jy$ (or 100 solar flux units) and a DL with $\ell = 100, R = 3km$, an ambient density $n = 10^9 cm^{-3}$, and a temperature of $10^6 K$, a brightness temperature of 3 $10^{22} K$ is required.

Such a brightness temperature is reached when radio waves of an initial temperature of $10^6 K$ are amplified with the growth rate of Eq. (21) over a distance

$$\Delta x \simeq 38c/\gamma_t \leq 10^4 c/\omega_{po} \simeq 3km \tag{24}$$

at an ambient plasma frequency of 300 MHz (density of $10^9 cm^{-3}$) and for $u^{3/2}\ell^{-2} \geq 10^{-2}$.

A theoretical upper limit to the brightness temperature of the radiation is obtained from the condition that the *radiation energy density should not exceed the electron beam energy density*

$$T_b^t < \frac{1}{2} \frac{n_f m v_f^2}{K} \cdot \frac{(2\pi)^3 R}{k_t^3 \ell \lambda_{Do}} \simeq 1.6\ 10^{25} T_6^{1/2} \frac{R}{3km} \frac{u^{1/2}}{\ell/100} K, \tag{25}$$

where we have used the vacuum dispersion relation $k_t = \omega_{po}/c$. One can further compare the radiation brightness temperature with the brightness temperature of the DL electric field

$$T_b^{DL} = \frac{E^2}{8\pi} \frac{\lambda_{Do} R^2}{K} \simeq 10^{23} \frac{(u/10)^2}{(\ell/100)^2} \left(\frac{R}{3km}\right)^2 T_6^{3/2} n_9^{1/2} K, \tag{26}$$

where we have used a wave vector volume for the DL of $(\lambda_{Do} R^2)^{-1}$. Spontaneous decay of radiation into DL field becomes only important for a radiation brightness temperature in excess of $(\omega^t/\omega^{DL}) T_b^{DL}$. So Eq. (26) does not lead to a stronger limitation for the brightness temperature of the radiation than Eq. (25).

The radiation is only amplified inside the DL and with maximum growth rate transverse to the DL electric field direction. The angular extent of the radiation depends therefore primarily on the thickness of the DL in the direction of the electric field with respect to its transverse dimension. For typical numbers used above one finds a *source which is radiating into a narrow equatorial double cone* of angular extent

$$\frac{L}{R} \simeq \frac{\ell\lambda_{Do}}{3km} \simeq 1 arc\ min. \tag{27}$$

https://doi.org/10.1017/S0074180900088276 Published online by Cambridge University Press

2.3 Relation with existing processes

The radiation treated above is a form of *Linear Accelerator Radiation* (Melrose, 1980; Krishan and Sivaram, 1983). We have shown that in the DL two coherent versions of this process exist, an antenna process and a maser. The maser mechanism is physically related to the processes of *scattering and double emission*. Further it is also instructive to compare our process with the *Free Electron Laser*, where low-frequency magnetic wiggles are upconverted by the (relativistic) electron beam into high- frequency radiation. In the case of the DL the low-frequency electric field takes over the role of the magnetic wiggles and the process operates already for nonrelativistic beams. In both cases *the energy of the radiation derives from the electron beam.*

3 Applications and discussion

We have investigated the possible existence of coherent radiation from DLs primarily because such structures may be the site of magnetic energy release *in solar and stellar flares* (Kuijpers, 1989a) and magnetic interactions between magnetospheres and accretion disks (Aly and Kuijpers, 1989). If this is true, coherent radiation from DLs would permit us to observe directly the primary sites of particle acceleration.

Here we propose that some of the *intense decimetric spikes* which are observed during solar flares (Kuijpers et al., 1981; Allaart and van Nieuwkoop, 1989) are the result of coherent radiation of DLs in the maser version treated in Section 2.2.

Further we propose that the *intense radio bursts of Type I* observed during solar noise storms are coherent radiation from transient DLs occurring in a coronal electric circuit triggered by reconnection. Evidence that the sources of Type I bursts are indeed the locations of electron acceleration comes from simultaneously observed Type III bursts which are commonly interpreted as electron beams and which seem to originate on the low-frequency side (higher altitude) of the Type I bursts (Leblanc, 1989).

One may ask how DLs are generated under cosmic conditions. In the laboratory and in numerical "experiments" one way to produce them is by driving current densities over the critical level (electron drift speed surpassing the electron thermal speed). In my opinion this is highly unlikely under cosmic conditions as this would require a constriction of the current cross-section by a factor

$$\frac{\omega_{co}}{\omega_{po}}\frac{c/\omega_{po}}{r/2}\frac{c}{v_{to}},$$

where ω_{co} is the electron cyclotron frequency, ω_{po} the electron plasmafrequency, v_{to} the electron thermal speed and r, the radius of the original current channel, is put equal to the scale height of the magnetic structure. However a different way to produce DLs in laboratory and numerical conditions is to excite a huge voltage spike locally. Under cosmic conditions this may be easy to establish by reconnection: If reconnection starts locally in a strongly sheared magnetic field system the global current system will oppose the forced change of the current path by *local inductive voltage spikes*. This may lead to the generation of a DL. The particle beams injected by this DL into the surrounding plasma may then cause anomalous resistivity which leads to enhanced dissipation, further voltage spikes and the generation of (transient) DLs. In this way local reconnection may trigger the global energy release in a cosmic electric circuit in the form of a multitude of transient DLs.

Coherent radiation from DLs in the maser version may also be relevant for some of the emission from *planetary magnetospheres*, complementary to the cyclotron maser. Direct observations of the terrestrial magnetosphere indeed show the existence of weak DLs or in any case of strong low-frequency electrostatic fields (Bostrom et al., 1989; Hultqvist, 1989). Although in our calculation we have used a model DL structure intense coherent radiation of both kinds (Sections 2.1 and 2.2) is generally expected if electron beams traverse electrostatic fields of sufficient intensity. The advantage of a DL is that it creates the beam itself and that its energy is externally maintained by the much larger energy of a current system.

Finally the antenna version of coherent DL emission may be relevant to as yet unexplained intense *microwave emission from a DL plasma in the laboratory* (Lindberg, 1988).

4 Conclusion

We have shown that nonrelativistic DLs emit intense radiation both of the antenna and of the maser type. The antenna version is suggested to be relevant to observed radiation in laboratory experiments on DLs. The maser version leads to observable emission in solar flares. Our calculations use several simplifying assumptions, notably a homogeneous electric structure with electric field parallel to the magnetic field, a strong potential difference inside the DL, weak turbulence methods, and finally dispersion relations for transverse waves which are strictly valid only in homogeneous plasmas and which neglect the magnetic field. We therefore consider these calculations as a first and promising step on the way to coherent radiation from realistic 3-D, magnetized, relativistic and non-relativistic, transient DL cavities.

Acknowledgement: I wish to thank Lennart Stenflo, Lennart Lindberg, Michael Raadu and Staffan Torvén for stimulating discussions on coherent radiation processes and DLs during my visit to Umea and Stockholm. The calculations were done during a stay at Meudon. In particular I thank Gérard Trottet, Monique Pick, Nicole Vilmer, Ludwig Klein, Bengt Pedersen, Yolande Leblanc, Philippe Zarka and other members of DASOP for the stimulating discussions and their hospitality.

References

Alfvén, H.: 1958, Tellus **10**, 104.

Alfvén, H.: 1981, Cosmic Plasma, D. Reidel Publ. Cy., Dordrecht, Netherlands.

Allaart, M.A.F. and van Nieuwkoop, J.: 1989, Solar Phys. submitted.

Aly, J.J. and Kuijpers: 1989, Astron. Astrophys., in press.

Bostrom, R., Holback, B., Holmgren, G., Koskinen, H.: 1988 Physica Scripta **39**, 782.

Hultqvist, B.: 1989, Phil. Trans. Royal Soc. London A **328**, 209.

Jacobsen, C. and Carlqvist, P.: 1964, Icarus **3**, 270.

Krishan, V. and Sivaram, C.: 1983, Solar Phys. **84**, 125.

Kuijpers, J.: 1989a, Solar Phys. **121**, 163.

Kuijpers, J.: 1989b, in Plasma Phenomema in the Solar Atmosphere, eds. M.A. Dubois and D. Grésillon, Cargèse, France.

Leblanc, Y.: 1989, in Plasma Phenomema in the Solar Atmosphere, eds. M.A. Dubois and D. Grésillon, Cargèse, France, in press.

Lindberg, L.: 1988, Astrophys. Space Sci. **144**, 3.

Melrose, D.B.: 1980, Plasma Astrophysics II, Gordon and Breach Publ., New York, p. 349.

Melrose, D.B.: 1986, Instabilities in Space and Laboratory Plasmas, Cambridge Univ. Press, Cambridge, U.K., Ch. 6.

Raadu, M.A.: 1984, Proc. 2nd Symp. on Plasma Double Layers, eds. R. Schrittwieser and G. Eder, Innsbruck, p. 3.

Raadu, M.A.: 1989, Physics Reports, **178**, 25.

Raadu, M.A. and Rasmussen. J.J.: 1988, Astrophys. Space Sci. **144**, 43.