

# Particle Acceleration by Induced Electric Fields in Course of Electric Current Oscillations in Coronal Magnetic Loops<sup>1</sup>

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**Abstract**— A mechanism for acceleration of electrons driven by the oscillations of the electric current in solar magnetic loops is considered. The magnetic loop is presented as an equivalent RLC-circuit with the electric current generated by convective motions in the photosphere. Eigen oscillations of the current in a loop induce the electric field directed along to the loop axis. It is shown that pulsating type III bursts and the sudden reductions that occur in the course of type IV continuum in solar flares provide evidence for acceleration and storage of energetic electrons in the coronal magnetic loops. Energization rate of electrons accelerated by sub-Dreicer electric field are determined. The energy spectra of fast electrons under both intermediate and strong regimes of pitch-angle diffusion are discussed. Two examples of the electron acceleration in the pulsating type III solar bursts are considered. We also discuss the efficiency of suggested mechanism as compared with the electron acceleration during 5-min photospheric oscillations and with the acceleration driven by the magnetic Rayleigh-Taylor instability.

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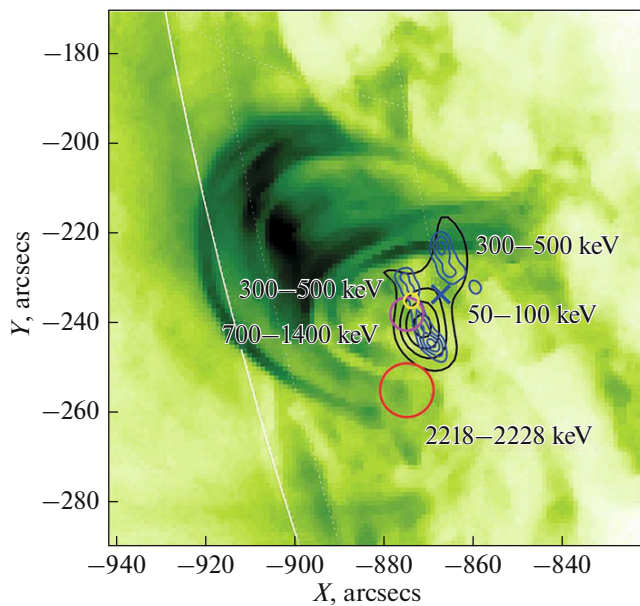
## 1. INTRODUCTION

There are evidences in favor of electron acceleration by electric field in solar flares. Nice example was presented by Lin et al. (2003) using RHESSI (gamma and hard X-rays) and TRACE (171 Å) observations of the flare of 2002 July 23. From Fig. 1 one can see that the centroid of the ion-produced 2.223 MeV neutron-capture line emission is located  $\sim 20'' \pm 6''$  away from the centroid of the 0.3–0.5 MeV and 0.7–1.4 MeV electron bremsstrahlung sources. It implies that the ions and electrons accelerate by a large-scale electric field because ions and electrons move in opposite directions and precipitate at different foot-points of the loop. But the question arises, what is the origin of the electric field? One of the possible origins of the electric field can be the magnetic reconnection. Craig and Litvinenko (2002) demonstrated that under the reconnection the strength of accelerating electric field  $E \sim B^2$ , and in the solar corona with magnetic field  $B \approx 100$  G the electrons may accelerate up to energies of several MeV. However there are no cusp structure of magnetic field in some flares and the coronal concept of the flare origin do not explain huge number of  $>20$  keV electrons ( $10^{39} - 10^{41}$ ) accelerated in the powerful flares (Miller et al., 1997). From the other side,

modern observations (Ji et al., 2012; Sharykin and Kosovichev, 2014) indicate that the main acceleration process occurs sometimes in the solar chromosphere. Zaitsev et al. (2015) have shown that the magnetic R-T instability near the chromospheric loop foot-points can be also the reason for induced electric field. Depending on the electric current in a loop this induced electric field can exceed the Dreicer field and thus the long-standing problem known as the “number problem” in the physics of solar flares (Hoyng, van Beek, and Brown, 1976) can be resolved.

Above mentioned mechanisms of the electron acceleration have the typical time scale from several tens of seconds to a few minutes and can not explain the multi-repeating acceleration processes as well as the long-lived events lasting for several tens of minutes. For example, observations performed by Aurass et al. (2003) revealed more than 100 broad-band radio pulsations in the flare of 25 October 1994 identified as the repeated type III bursts and interpreted by Zlotnik et al. (2003) as the numerous electron beam injections into an asymmetric magnetic loop (Fig. 2). Observations of repetitive particle accelerations in HXR, microwave, and type III radio bursts with a period of about 3 minutes in the flare on 23–24 September 2011 were presented by Kumar et al. (2016). Quasi-periodic bunches of type III bursts with period of 270 s and with the positive frequency drift cover the range 0.7–3 GHz

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**Fig. 1.** Flare on July 23, 2002 obtained with the space observatories. *RHESSI* image is superposed on a *TRACE* map of the loops taken 90 minutes after the flare. White line indicates the solar limb. The sources of hard X-ray emission 50–1400 keV (high resolution) generated by accelerated electrons in the loop foot-points and gamma radiation in the 2.2 MeV lines (low resolution) generated by the accelerated ions are spatially separated. Ions and electrons move in opposite directions and bombard thick targets in different foot points of the coronal magnetic loops (Lin et al., 2003).

in the flare on 2012 July 19 were investigated by Huang et al. (2016).

Periodical type III bursts promote accumulation of fast particles in coronal loops. Zlotnik et al. (2003) showed that the broad band radio pulsations in the event of 25 October 1994 resulted from periodically

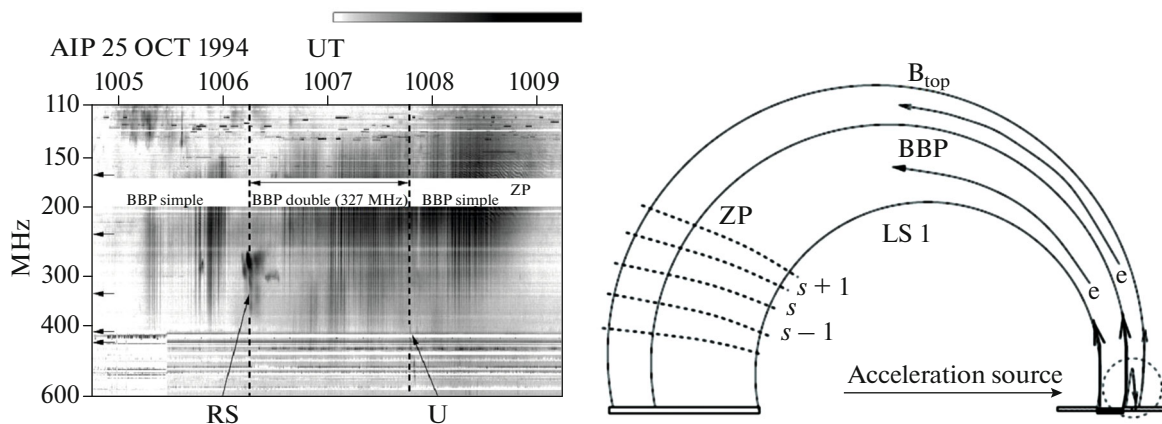
repeated injections of fast electrons into the magnetic trap. Zebra patterns in the event on 25 October, 1994 arose 2.5 min after the broad band pulsations. If the sources of both fine structures are located in the same magnetic loop, this time delay means that the Zebra pattern starts after about a hundred injections of fast electrons into the magnetic trap. This is because the instability threshold for the wave excitation at the double plasma resonance needed for Zebra pattern exceeds the threshold of beam instability which causes broad band pulsations in type III bursts.

Sudden reductions in the flaring radio emission also indicate on the storage of energetic electrons in the coronal loops. If we assume that the continuum emission of type IV is caused by the loss-cone instability in the loop (a magnetic trap), the sudden reductions can occur as a result of injection of accelerated electrons into the trap (Zaitsev and Stepanov, 1975; Benz and Kuijpers, 1976). A part of the injected particles fills the loss-cone and quenches the loss-cone instability as the cause of the Type IV emission. Wave-particle interaction changes the pitch-angle distribution of injected fast electrons; therefore, the number of energetic particles trapped in a loop grows.

The idea of the acceleration and storage of energetic electrons in solar magnetic loops by the electric current oscillations was suggested by Zlotnik et al. (2003). Here, we will develop the idea of acceleration and accumulation of the magnetic loops by energetic particles in the course of the electric current oscillations.

## 2. EXCITATION OF INDUCED ELECTRIC FIELD BY CURRENT OSCILLATIONS

Convective motions of the photospheric plasma interacting with the magnetic field near the loop foot-points generate electric current inside the loop. This current flows through the coronal part of the loop



**Fig. 2.** Left: The 40–800 MHz spectrum of the 25 October 1994 event with > 100 type III bursts. BBP – broad band pulsations; (Aurass et al., 2003). Right: A source model. The loop LS1 is the main source of BBP. Fast electron streams exciting BBP are injected from the right footpoint (Zlotnik et al., 2003).

from one of its footpoints to the other and closes near or inside the photosphere, where the conductivity is isotropic. Hence, the current-carrying magnetic loop can be presented as an equivalent electric circuit ( $RLC$ ), whose eigen-frequency depends on the magnitude of the constant component of the electric current  $I_0$ , number density  $n$ , loop radius  $r_0$ , and the length  $l$  of the coronal part of the loop (see for example, Stepanov, Zaitsev, and Nakariakov, 2012):

$$v_{RLC} \approx \frac{1}{2\pi\sqrt{2\pi\Lambda}} \frac{I_0}{cr_0^2\sqrt{nm_i}}, \quad \Lambda = \ln\left(\frac{4l}{\pi r_0}\right) - \frac{7}{4}. \quad (1)$$

Electric current oscillations in the equivalent  $RLC$ -circuit are connected with the oscillations of the azimuthal component of the magnetic field in the loop

$$B_\phi(r, t) = 2r \frac{I_z(t)}{cr_0^2}. \quad (2)$$

These oscillations, in turn, in accordance with the Maxwell-Faraday equation

$$\text{rot} \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}_\phi}{\partial t}, \quad (3)$$

lead to the generation of an electric field directed along the flux tube axis. This electric field is parallel to the component of the magnetic field  $B_z$ , therefore, it should efficiently accelerate charged particles. Assuming  $I_z(t) = I_0 + \Delta I \sin(2\pi v_{RLC}t)$  and averaging over the half-period of the oscillations with taking into account Eqs. (2) and (3) we found the relations of the electric field with the loop radius  $E_z(r)$ , the maximum value  $E_{\max}$  reached near the loop surface, and the mean value along the loop radius  $\bar{E}$ :

$$\begin{aligned} E_z(r) &= \frac{4v_{RLC}I_0\Delta I}{c^2} \frac{I_0}{I_0} \left(\frac{r}{r_0}\right)^2, \\ E_{\max} &= \frac{4v_{RLC}I_0\Delta I}{c^2} \frac{I_0}{I_0}, \\ \bar{E} &= \frac{4}{3} \frac{v_{RLC}I_0\Delta I}{c^2} \frac{I_0}{I_0}. \end{aligned} \quad (4)$$

From Eq. (4) one can see that the accelerating electric field depends on the amplitude of electric current oscillations  $\Delta I$ . In the self-consistent equation for equivalent  $RLC$  circuit, the resistance and the capacitance turn out to be dependent on the electric current (Khodachenko et al., 2009):

$$\begin{aligned} \frac{1}{c^2} L \frac{d^2 y}{dt^2} + \left[ R(I_0)(1+y)^2 - \frac{|V_r|l_1}{c^2 r_1} \right] \frac{dy}{dt} \\ + \frac{1}{C(I_0)} \left( 1 + \frac{3}{2}y + \frac{1}{2}y^2 \right) y = 0. \end{aligned} \quad (5)$$

Here,  $y = (I - I_0)/I_0$ ,  $L$  is the loop inductance. If the current-carrying loop is approximated by a circular wire with a length  $l$  and a small radius  $r_0 \ll l$ , for the

inductance one can use the well-known expression (Landau and Lifshitz, 1984):

$$L \approx 2l\Lambda = 2l \left( \ln \frac{4l}{\pi r_0} - \frac{7}{4} \right), \quad (6)$$

where  $C(I_0)$  is the capacitance depending on the current flowing in the loop:

$$\frac{1}{C(I_0)} \approx \frac{I_0^2 l}{c^4 n m_i \pi r_0^4} (1 + b^{-2}). \quad (7)$$

Parameter  $b$  is determined by the components of the magnetic field and the plasma gas pressure both at the axis of the flux tube and outside it:

$$b = \frac{B_{\phi 0}(r_0)}{B_{z 0}(r_0) - B_{z 0}(0)} \approx 6 \frac{B_{\phi 0}(r_0) B_{z 0}(0)}{8\pi [p(\infty) - p(0)]}. \quad (8)$$

Here,  $B_{\phi 0}$  and  $B_{z 0}$  are azimuthal and axial components of the magnetic field of the magnetic loop. Effective resistance of a loop is determined as

$$R(I_0) \approx \frac{F_1^2 I_0^2 l_1}{(2 - F_1) c^4 n_i m_i v_{ia} \pi r_1^4} (1 + b^{-2}), \quad (9)$$

where  $l_1, r_1, n_i, F_1$  are the scale of the loop along the height, loop radius, plasma density, and relative density of neutral particles in the area of the photospheric  $emf$ , respectively, and  $v_{ia}$  is the frequency of electron-atom collisions. The right term in the square brackets in Eq. (5) denotes  $emf$  in the foot point of the loop and can be considered as a negative resistance;  $V_r$  is the radial component of the convergent flow of the photosphere matter in the loop foot-points. The horizontal flow of plasma toward the loop is driven by the photosphere convection that occurs, for example, when the footpoints of the loop are in the node of several supergranulation cells. The footpoints of the loop as an equivalent circuit primarily contribute to the whole loop resistance due to their comparatively low conductivity determined by ion-atom collisions (the so-called Cowling conductivity). In the coronal part of the loop,  $b^2 \gg 1$ , therefore, for the frequency of  $RLC$  oscillations we obtain Eq. (1).

Eq. (5) suggests that electric field oscillations must be in-phase at all points of the loop (lumped circuit approach). On the other hand, the current variations propagate along the loop with the Alfvén velocity. Therefore, for the in-phase condition, the Alfvén time  $\tau_A = l/V_A$  must be less than the period of oscillations  $T_{RLC} = 1/v_{RLC}$ . Since the Biot-Savart law suggests  $I_0 \approx cr_0 B_{\phi 0}(r_0)/2$ , the twisting of the loop magnetic field should be low (Stepanov, Zaitsev, and Nakariakov, 2012):

$$\frac{B_{\phi 0}(r_0)}{B_{z 0}(0)} < 2\pi\sqrt{2\Lambda} \frac{r_0}{l}. \quad (10)$$

Eq. (1) corresponds to the frequency of Alfvén oscillations in a coronal magnetic loop with a wave vector  $|\vec{k}| \approx r_0^{-1}$  directed almost perpendicular to the axis of the magnetic flux tube at an angle  $\cos \theta \approx (B_\phi/B_z)$  to the magnetic field  $\vec{B}$ . In this case, the frequency of Alfvén oscillations is

$$v_A = \frac{1}{2\pi} k V_A \cos \theta \approx \frac{1}{2\pi r_0} \frac{\sqrt{B_z^2 + B_\phi^2} B_\phi}{\sqrt{4\pi n m_i} B_z} \quad (11)$$

$$\approx \frac{1}{2\pi\sqrt{\pi} c r_0^2 \sqrt{n m_i}},$$

which, up to a coefficient of the order of unity, coincides with the Eq. (1). Here, we took into account that  $I_0 \approx B_\phi c r_0/2$ .

From Eq. (5) it also follows that the excitation of the oscillations of the loop as an equivalent RLC circuit occurs if the negative resistance of the photosphere *emf* exceeds the loop resistance:

$$R(I_0) \leq |V_r| l_i / (r_i c^2). \quad (12)$$

The  $Q$ -factor of the equivalent electric circuit is

$$Q = \frac{1}{cR} \sqrt{L/C}, \quad (13)$$

and can be very high because of the high loop inductance and comparatively low circuit resistance. Introducing the dimensionless time  $\tau = 2\pi\nu_{RLC} t$  we can write Eq. (5) in the form

$$\frac{d^2 y}{d\tau^2} - \varepsilon (\delta - 2y - y^2) \frac{dy}{d\tau} + \left(1 + \frac{3}{2}y + \frac{1}{2}y^2\right) y = 0. \quad (14)$$

Here,  $\varepsilon = 1/Q$ ,  $\delta = \left[|V_r| l_i / c^2 r_i R(I_0)\right] - 1$ . The small parameter  $\varepsilon \ll 1$  in Equation (14) makes it possible to apply the Van der Pol method to solve the Equation; under the steady-state condition, the solution has the form (Zaitsev, Stepanov, and Kaufmann, 2014)

$$y = \frac{I(t) - I_0}{I_0} = 2\sqrt{\delta} \cos \left[ 2\pi\nu_{RLC} \left(1 + \frac{3}{4}\delta\right) t \right]. \quad (15)$$

Thus, the nonlinearity in Eq. (5) leads to the establishment of finite amplitude of the oscillations, as well as to a small shift in the oscillation frequency in comparison with the linear case. The magnitude of the electric current oscillations in the steady-state regime is determined by the excess of the negative resistance of the photospheric *emf* located in the coronal loop footpoints over the active resistance of the equivalent electric circuit. Kaufmann et al. (2009) observed the high- $Q$  pulsations of sub-THz radiation in the solar flare on November 4, 2003 which was interpreted by Zaitsev, Stepanov, and Kaufmann (2014) as the effect of modulation of the emission volume by eigen-oscillations of the current-carrying magnetic loop. These authors showed that the amplitude of pulsations of sub-THz

radiation is proportional to the amplitude of the electric current variations in the coronal magnetic loop, since the modulation of the flux of energetic electrons accelerated by the induction electric field (Eq. 4) depends on the amplitude of current variations. This dependence makes it possible to determine for the flare on November 4, 2003 the amplitude of the electric current oscillations in the coronal magnetic loop using the data on high- $Q$  pulsations of sub-THz radiation,

$$\Delta I/I_0 = 2\sqrt{\delta} \approx (5-8) \times 10^{-2}, \quad (16)$$

as well as the magnitude of the accelerating electric field,

$$E_z(r) \approx (0.2-0.3) \frac{V_{RLC} I_0}{c^2} \left(\frac{r}{r_0}\right)^2. \quad (17)$$

As it follows also from Eqs. (1) and (4)  $E_z(r)$  quadratically depends on the electric current in the loop and increases from the loop axis to the loop surface.

### 3. DENSITY AND ENERGY SPECTRUM OF ACCELERATED ELECTRONS

The induced electric field (Eqs. 2 and 17) accelerates some part of electron population with the velocity exceeding  $V > (E_D/E_z)^{1/2} V_{Te}$ , where  $V_{Te} = (k_B T/m)^{1/2}$  is the electron thermal velocity,  $k_B$  is the Boltzmann constant,  $E_D = e\Lambda_c \omega_p^2 / V_{Te}^2$  the Dreicer field,  $\Lambda_c$  the Coulomb logarithm,  $\omega_p$  the Langmuir frequency. In the case of sub-Dreicer field, *i.e.*  $x = E_D/E_z \gg 1$ , the kinetic theory yields the number of runaway electrons per second accelerated in DC-electric field (Knoepfel and Spong, 1979):

$$\dot{n}_S = 0.35 n v_{ei} V_a x^{3/8} \exp\left(-\sqrt{2x} - \frac{x}{4}\right), \quad (18)$$

where  $v_{ei} = 5.5 n \Lambda_c / T^{3/2}$  is the effective frequency of electron-ion collisions,  $n$  the thermal electron number density,  $T$  the plasma temperature,  $V_a$  the volume of the acceleration region. From Eqs. (17) and (18) it follows that in the case of sub-Dreicer electric field, the flux of accelerated electrons reaches its maximum near the surface of the magnetic flux tube in a layer with the thickness

$$\Delta r \approx \frac{r_0}{\sqrt{2x_m} + \frac{x_m}{2}}, \quad (19)$$

where  $x_m = E_D/E_{\max}$ .

The total number of electrons accelerated in a trap per unit time can approximately be estimated assuming their isotropic distribution over the pitch-angles

$$\dot{N}_s \approx 2\dot{n}_s(x_m)\sigma S_0 l \frac{\Delta r}{r_0}, \quad (20)$$

where  $\sigma = B_0/B_{\text{top}}$  is the mirror ratio of the trap with the magnetic field  $B_0$  at the footpoint and  $B_{\text{top}}$  at the loop top, and  $S_0$  is the foot cross-section area.

The regimes of pitch-angle diffusion of accelerated particles into the loss cone depend on the ratio of the following time scales (Bespalov and Trakhtengerts, 1986) the diffusion time  $\tau_D$ , which corresponds to the mean time for the change in the pitch-angle of a fast electron at  $\pi/2$  due to the collisions with particles or the wave-particle interaction, the time of filling the loss cone  $\tau_D/\sigma$ , and the time of flight when escaping into the loss cone,  $\tau_0 = l/2V$ . In the case of intermediate diffusion, when the condition  $\tau_0 < \tau_D < \sigma\tau_0$  is satisfied, the loss cone is filled faster than the particles escape into the loss cone. In this case, the lifetime of a particle in a magnetic trap is of the order of magnitude  $\tau_i \approx \sigma\tau_0$ , and the particle loss rate is  $\dot{N}_{\text{loss}} \approx N_s/\tau_i \approx 2n_s\sigma l S_0(\Delta r/r)/\tau_i$ . Then from the condition  $\dot{N}_s \approx \dot{N}_{\text{loss}}$  we can estimate of the density of fast electrons with the energy  $\varepsilon$  in the acceleration region (in a shell of the thickness  $\Delta r$ ):

$$n_s(\varepsilon) \approx \frac{1}{2\varepsilon_m} \dot{n}_s(x_m) \frac{\sigma l}{\sqrt{2\varepsilon/m}}. \quad (21)$$

The energy dispersion of accelerated electrons is determined by the interval

$$k_B T \left( \frac{E_D}{E_{\text{max}}} \right) \leq \varepsilon \leq \varepsilon_m = eE_{\text{max}}l, \quad (22)$$

within which the velocity of formation of runaway electrons (Eq. 18) does not depend on their energy. From Eq. (21) we obtain the total density of accelerated electrons:

$$n_s \approx \int n_s(\varepsilon)d\varepsilon \approx \frac{\dot{n}_s(x_m)\sigma l}{\sqrt{2\varepsilon_m/m}}. \quad (23)$$

In the case of strong diffusion, when the condition  $\tau_D < \tau_0$  is satisfied, the average time of pitch-angle diffusion is less than the time for free escape of the particle into the loss cone. It means that during the traverse of the loop length, the particle repeatedly changes the direction of motion and the lifetime of the particle in the trap increases. For example, in the case of diffusion on whistlers, the lifetime of fast electrons is determined by the formula (Bespalov and Trakhtengerts, 1986)

$$\tau_i \approx \frac{m_e}{m_i} \left( \frac{V}{V_A} \right)^2 \sigma \tau_0. \quad (24)$$

In this case the density of accelerated electrons with the energy  $\varepsilon$  is

$$n_s(\varepsilon) \approx \frac{1}{2\varepsilon_m} \dot{n}_s(x_m) \frac{\sigma l}{V_A^2} \frac{m}{m_i} \sqrt{\frac{2\varepsilon}{m}}, \quad (25)$$

and the total density is as follow

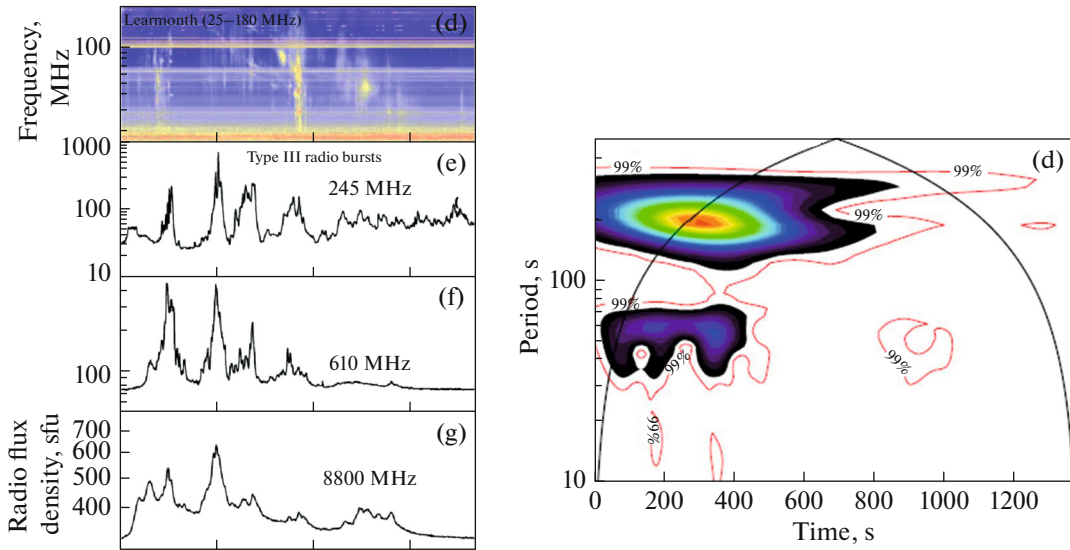
$$n_s \approx \int n_s(\varepsilon)d\varepsilon \approx \frac{\dot{n}_s(x_m)\sigma l}{3V_A^2} \frac{m}{m_i} \sqrt{\frac{2\varepsilon_m}{m}}. \quad (26)$$

From Eqs. (21) and (25) one can see that in the case of intermediate diffusion, the spectrum of accelerated particles decreases with increasing of the energy. In the case of the strong diffusion regime, the spectrum increases with the energy increasing. It should be noted that formulas (23) and (26) determine the density of accelerated particles in a loop – a magnetic trap, with a uniform distribution of the background plasma density. If the latter depends on the coordinate along the loop, then the parameter  $x_m$  also varies with the coordinate and this lead to a change in the spectrum of the accelerated electrons as compared to Eqs. (21) and (25). Therefore, formulas (23) and (26) are determined correctly the total density of accelerated electrons in the energy interval (22), while the real power-law energy spectra can depend both on the inhomogeneity of the plasma inside the loop and on a number of other factors, for example, on Coulomb losses, wave-particle interaction, etc.

#### 4. TWO EXAMPLES OF PULSING TYPE III BURSTS

In order to illustrate of proposed model let us consider two solar events with the quasi-periodic type III bursts and estimate the acceleration rate and the energy of accelerated electrons.

(i) Flare on 23–24 September 2011 studied by Kumar et al. (2016) reveals the 180-s quasi-periodic pulsations observed in the meter and decimeter radio wavelengths as the repetitive bunches of type III bursts with negative frequency drift. It means that the acceleration region is located below the level at which the Langmuir frequency  $\omega_p/2\pi \approx 600$  MHz. Assuming for the coronal loop  $n = 10^{10} \text{ cm}^{-3}$ ,  $T = 10^6 \text{ K}$ ,  $l = 10^{10} \text{ cm}$ ,  $r_0 = 10^8 \text{ cm}$ , and the electric current  $I_0 = 2 \times 10^9 \text{ A}$ , from Eq. (1) we can estimate the eigen-frequency of the current-carrying loop,  $\nu_{RLC} \approx 5.54 \times 10^{-3} \text{ Hz}$ . Thus, the period of RLC-oscillations is 180 s, equal to the pulsation period observed in the repetitive type III bursts event on September 23–24, 2011. Taking a small modulation magnitude,  $\Delta I/I_0 = 10^{-3}$ , from Eq. (2) we obtain the mean value of the electric field along the loop radius:  $\bar{E} \approx 1.5 \times 10^{-5} \text{ V cm}^{-1}$ . Electrons can be accelerated by this electric field at the distance  $\Delta l = 2 \times 10^9 \text{ cm}$  up to the energy  $\varepsilon \approx 30 \text{ keV}$ , the typical energy of electrons responsible for type III bursts. The Dreicer field is  $E_D = 6 \times 10^{-8} n/T = 6 \times 10^{-4} \text{ V cm}^{-1}$ , and



**Fig. 3.** Left: (d) dynamic radio spectrum in 25–180 MHz frequency, showing bunches of type III radio bursts. (e)–(g) RSTN radio flux density profiles in 245, 610, and 8800 MHz frequencies from the Learmonth solar observatory. Right: Wavelet power spectra for RSTN 15.4 GHz signal. The start time is 2348 UT on 23 September 2011 (Kumar et al., 2016).

the ratio  $x = E_D/\bar{E} \approx 40 \gg 1$ . Supposing the acceleration volume  $V_a = 6 \times 10^{25} \text{ cm}^3$ , from Eq. (18) we estimate the energization rate  $\dot{n}_s \approx 3 \times 10^{30} \text{ el s}^{-1}$ , which is about four times higher than in the pulsating type III event on 25 October 1994 (Zlotnik et al., 2003), but 4–6 orders of magnitude lower than the rate of electron acceleration in large flares.

For this estimate, we used the mean value of the induced electric field along the loop radius  $\bar{E}$ . The maximum value  $E_{\text{max}}$  (Eq. 2) is reached near the loop surface and the maximum fluxes of accelerated electrons occur near the surface of the magnetic flux tube in a layer with the thickness (Eq. 19)  $\Delta r = 0.14r_0 \approx 3 \times 10^7 \text{ cm}$ . As the induced electric field quadratically depends on the electric current, a part of the electron population accelerated near the loop surface can reach the energy of about 1 MeV. Moreover, some fraction of energetic electrons can propagate along the open magnetic field lines as type III bursts due to diffusion across magnetic field or due to the coalescence between closed and open magnetic field lines. An important peculiarity of RLC-oscillation of the coronal magnetic loops is a high  $Q$ -factor (Eq. 13), which is consistent with the idea of pumping of magnetic loops by the energetic electrons. Indeed, using Eqs. (6) and (7) and taking the threshold value of  $R(I_0)$  in Eq. (9) to be  $|V_1| = 5 \times 10^4 \text{ cm s}^{-1}$ ,  $l_1 = 10^8 \text{ cm}$ ,  $r_1 = 5 \times 10^7 \text{ cm}$ , for the plasma loop we obtain the  $Q$ -factor value,  $Q \approx 10^4 \gg 1$ .

(ii) In the flare on July 19, 2012 (Huang et al., 2016), the set of type III bursts displays in-phase oscillations at 270 s and a positive frequency drift within the range of 0.7–3 GHz (Fig. 4) The positive frequency

drift indicates that the region of electron acceleration is located near the loop top and the beams of energetic electrons are directed downward. The Dreicer electric field reaches its minimum in this region of a loop, which favors the acceleration process. Taking  $n = 6 \times 10^9 \text{ cm}^{-3}$ ,  $T = 10^6 \text{ K}$ ,  $l = 1.5 \times 10^{10} \text{ cm}$ ,  $r_0 = 2 \times 10^8 \text{ cm}$ , and  $I_0 = 4 \times 10^9 \text{ A}$ , from Eq. (1) we obtain the eigen-frequency of RLC-oscillations,  $\nu_{RLC} \approx 3.76 \times 10^{-3} \text{ Hz}$ . The period of oscillations (266 s) is close to that observed.

The mean value of the electric field for  $\Delta I/I_0 = 10^{-3}$  is  $\bar{E} \approx 2 \times 10^{-5} \text{ V cm}^{-1}$ , and  $E_D \approx 4 \times 10^{-4} \text{ V cm}^{-1}$ , thus  $x \approx 20$ . Assuming the acceleration volume  $V_a = 10^{26} \text{ cm}^3$  from Equation (18) we can derive an effective energization rate  $\dot{N}_s \approx 3 \times 10^{33} \text{ el s}^{-1}$ . This is compatible with the energization rate obtained from the RHESSI data (Holman et al., 2011). The energy gain for  $\bar{E} \approx 2 \times 10^{-5} \text{ V cm}^{-1}$  at the distance  $\Delta l = 10^9 \text{ cm}$  is approximately  $\varepsilon \approx 20 \text{ keV}$ .

Huang et al. (2016) suggested that the acceleration site of electrons producing type III bursts is located above the loop top. In our interpretation, the acceleration site is located in the place with the minimum of the Dreicer field, *i.e.* at the loop top. Note also that the lumped circuit approach applied here and suggesting  $\nu_{RLC} < V_A/l \approx 10^{-2} \text{ Hz}$  is satisfied for both of the considered events.

## 5. DISCUSSION

In this Section we will discuss the efficiency of the acceleration mechanism suggested here in comparison

with the electron acceleration driven by 5 min oscillations of the photosphere (Zaitsev and Kislyakov, 2006), and the acceleration due to the Rayleigh-Taylor instability (Zaitsev and Stepanov, 2015).

### 5.1. Acceleration Driven by 5-min Oscillations of the Photosphere Convection

It is well known that the photosphere convection displays a wide spectrum of velocity oscillations with the maximum of about five minutes. These oscillations have the acoustic origin, however, they do not penetrate into the chromosphere and corona, being reflected from the region of temperature minimum. Thereby, the 5-min oscillations cannot directly affect the entire coronal magnetic loop. Nevertheless, acoustic oscillations modulate the electric currents flowing in the coronal loops due to the “intermeshing” of the convective plasma flow and the magnetic field in the loop foot-points. The modulation of the electric current leads, in turn, to periodic changes in the electric field directed along the axis of the flux tube.

Let us assume that the radial component of the photosphere convection velocity is modulated by 5-min photosphere oscillations according to the law

$$|V_r| = V_0 + V_- \sin(\omega_5 t), \quad (27)$$

where  $V_0 \gg V_-$ . Slow variations of the current in the loop on a timescale longer than the period of eigen-oscillations of the loop as the RLC circuit are described by the equation (Stepanov, Zaitsev, and Nakariakov, 2012):

$$\frac{L}{c^2} \frac{\partial I_z}{\partial t} + RI_z = \Xi(I_z). \quad (28)$$

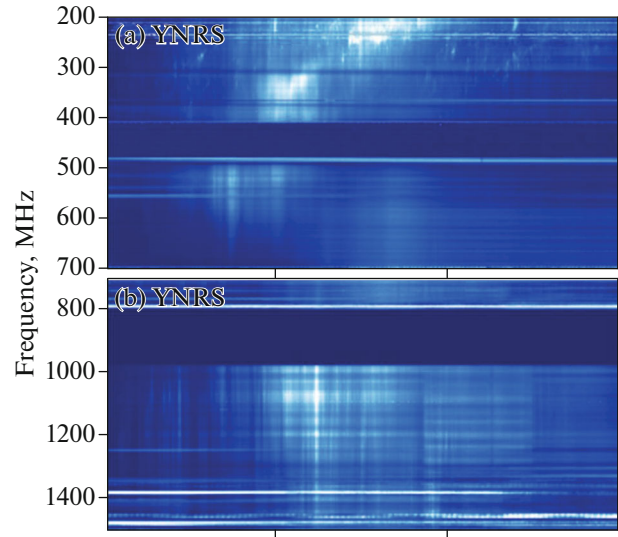
Here *emf*  $\Xi(I_z)$  generating the current flow along the loop is connected with the velocity of photospheric convection through the relation

$$\Xi(I_z) = \frac{V_r B_\phi h}{c} \cong \frac{|V_r| l_1}{r_1 c^2} I_z. \quad (29)$$

The height interval extends from the lower layers of the photosphere to the transition region between the photosphere and the chromosphere and for the solar atmosphere is roughly equal to  $l_1 = 500\text{--}1000$  km. The oscillations in the velocity of photosphere convection  $V_-$  are connected with the oscillations of the electric current in the loop,  $I_z = I_0 + I_-$ , and for  $I_-$  we have the equation

$$\frac{\partial I_-}{\partial t} + \frac{2R(I_0)c^2}{L} I_- = \frac{V_- l_1 I_0}{r_1 L} \sin \omega_5 t. \quad (30)$$

This equation has a forced periodic solution for a current that varies with the frequency of the photospheric oscillations  $\omega_5$  and displays the relative amplitude of oscillations  $I_-^m/I_0 \approx l_1 V_- / \omega_5 L r_1$  (Zaitsev and Kislyakov,



**Fig. 4.** Decimetric and metric dynamic spectra of the flare July 19, 2012 with a bunches of pulsing type III bursts (Huang et al., 2016). Time interval covers 30 min.

2006). By analogy with Equation (2), we obtain the expressions for accelerating electric field

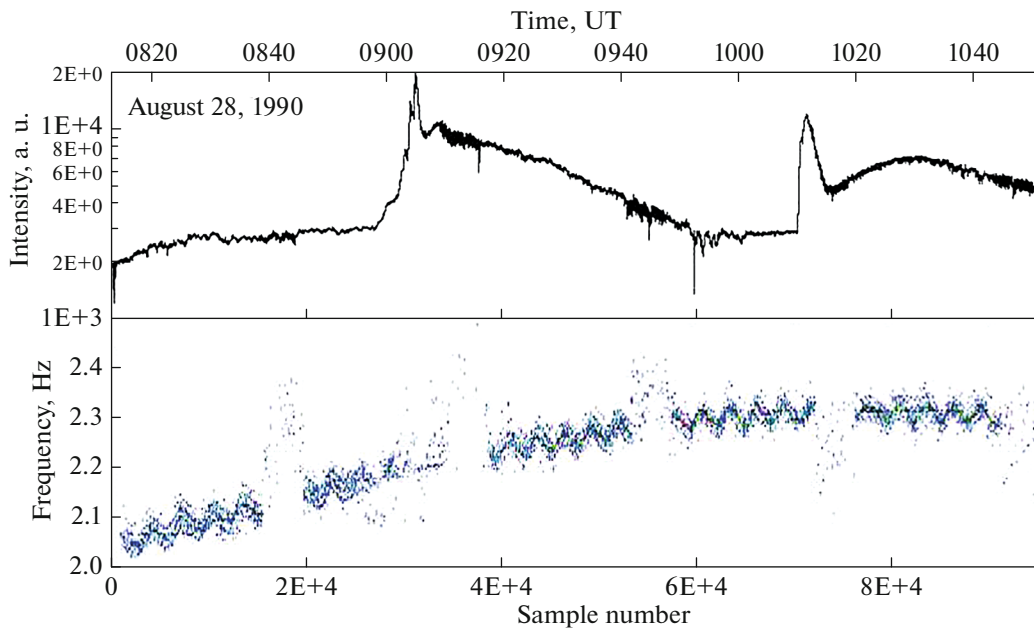
$$E_z(r) = \frac{4v_5 I_0 I_-^m}{c^2 I_0} \left( \frac{r}{r_0} \right)^2, \quad E_{\max} = \frac{4v_5 I_0 I_-^m}{c^2 I_0}, \quad (31)$$

$$\bar{E} = \frac{4}{3} \frac{v_5 I_0 I_-^m}{c^2 I_0},$$

where  $v_5 = \omega_5 / 2\pi$ . The upper panel in Fig. 5 shows two solar radio bursts observed at 37 GHz with Metsähovi radio telescope on 28 August 1990, and the lower panel presents the modulation of eigen-oscillations of the loop as an equivalent electric circuit by 5-min photospheric oscillations obtained with the Wigner-Wille transform. The latter transform displays appreciably higher frequency-time resolution. It can be seen that the relative variations of the frequency of the eigen-oscillations of the RLC-circuit caused by the oscillations of the electric current have a value of  $I_-^m/I_0 \approx \Delta v_{LRC} / v_{LRC} \approx (1-2) \times 10^{-2}$ , which coincides with the above estimate for pulsations of sub-THz radiation (Zaitsev, Stepanov, and Kaufmann, 2014). Comparing Eqs. (4) and (31) we see that 5-min photospheric oscillations induce approximately the same electric fields in a coronal magnetic loop as the eigen-oscillations of the loop do.

### 5.2. Acceleration Driven by the Rayleigh-Taylor Instability

It is also possible that intense electric fields in the magnetic loops are generated as a result of the magnetic Rayleigh-Taylor (R-T) instability in the chromospheric footpoints of the loops. This instability leads to



**Fig. 5.** Modulation of microwave radiation at 37 GHz by 5 min photosphere oscillations in the flare on 28 August 1990 (Zaitsev et al., 2003). Upper panel shows light curve. Lower panel presents the dynamic spectrum of low-frequency modulation of the microwave emission flux obtained by the Wigner-Wille transform.

penetration of external chromospheric plasma with  $\beta = 8\pi p/B^2 \leq 1$  into a magnetic loop, where  $\beta \ll 1$ . Here  $p$  is the plasma gas pressure. As a result, the magnetic field of the loop is deformed and amplifies the electrical current. Directly in the footpoint of the magnetic loop, where the R-T instability develops, the induced electric field is perpendicular to the magnetic field,  $\vec{E} = -(1/c)[\vec{V} \times \vec{B}]$ , hence this electric field is not able to accelerate particles. However, for the time  $\tau_A \approx \Delta l/V_A \approx 5-25$  s, where  $\Delta l \approx (1-5) \times 10^7$  cm is the scale of a plasma tongue invading the magnetic loop, a pulse of the tension of the magnetic field  $B_\phi^2(r, t)/8\pi$  leaves the R-T instability region along the flux tube axis in the form of a pulse of a longitudinal electric current. (Zaitsev et al., 2016) have shown that if the electric current is sufficiently strong,  $B_\phi^2 \gg 8\pi p$ , an induced electric field directed along the magnetic field of a loop  $B_{z0}$  appears:

$$\frac{\partial E_z}{\partial r} = -\frac{1}{c} \frac{B_\phi^2}{4\pi\rho V_A^2} \frac{\partial B_\phi}{\partial t}. \quad (32)$$

This is due to the fact that for  $B_\phi^2 \gg 8\pi p$  the perturbations of the magnetic field are not compensated by the gas pressure gradient as it occurs in the linear Alfvén pulse. The perturbations of the velocity along the flux tube radius and along  $B_{z0}$  appear which results in generation of electric field component along the flux tube axis which is non-linear with respect to the magnetic field component  $B_\phi$ . The maximum of the is

reached at  $r = r_0$  and its ratio to the Dreicer field is equal to

$$\frac{E_{zm}}{E_D} = 2.2 \times 10^8 \frac{TI_0^3(A)}{r_0^2 B_{z0} n^{3/2} \Delta \xi}, \quad (33)$$

where  $\Delta \xi$  is the scale of electric current pulse propagating in a coronal magnetic loop along  $B_{z0}$ . Eq. (33) suggest that the induced electric field  $E_z$ , depends on the current as  $E_z \propto I^3$  and can exceed the Dreicer field for  $I_0 > 10^{10}$  A (Zaitsev et al., 2016). Hence, acceleration driven by the R-T instability is the most efficient mechanism because at sufficiently high values of the electric current the inductive electric field in the chromosphere can be of the order or even more the Dreicer field, greatly increasing the number of accelerated electrons. However it requires quite high electric currents,  $I_0 > 10^{10}$  A, and the typical time of the R-T acceleration can not explain the long-lived radio events.

## 6. CONCLUSIONS

Several models have been suggested to interpret quasi-periodic acceleration of energetic particles in the solar flares. The MHD oscillation of flare loops could be a possible option to explain the features under question. However, the sausage and kink MHD modes are not capable to provide synchronous pulsations in a wide frequency interval and have quite low  $Q$ -factor. Therefore, they are inappropriate as a cause for the acceleration in the coronal loops.



Kuijpers (1978) has shown that plasmaturbulence in solar flares may lead to the periodic acceleration of charged particles. His proposal is that the periodicity of the acceleration process is due to the interaction of ion-sound and Langmuir waves described by the Lotka-Volterra equations. But the model of Kuijpers has the stochastic origin, and deals not to the induced electric field. In addition it can't explain quasi-periodical structures with the period of about several minutes. Quasi-periodic acceleration may be associated with the bursty regime of spontaneous magnetic reconnection, such as that in the case of tearing of the current sheet associated with the formation of multiple plasmoids (Kliem et al., 2000). Kumar et al. (2016) also suggest that the periodic reconnection at a magnetic null point most likely causes the repetitive particle acceleration. Nevertheless, all the reconnection regimes cannot explain the high  $Q$ -factor of the observed acceleration. Indeed, simulations performed by Tajima et al. (1987) show that the energy of electrostatic and inductive fields decreases by an order of magnitude already after the first few oscillations.

The mechanism of acceleration and storage of electrons driven by the electric current oscillations in a loop as an equivalent RLC-circuit was first suggested in the paper of Zlotnik et al. (2003). Here, we develop this idea but for the periodic bunches of type III bursts. By this the individual type III bursts forming the bunches can be triggered by the bursty regime of magnetic reconnection in the loop structure, the thin flux tubes with cross-section area of about  $10^{14}$ – $10^{15}$  cm<sup>2</sup>, and/or by the sausage mode pulsations of these fine structures. The proposed model is similar to that of electron acceleration initiated by the 5 min velocity oscillations of the photosphere (Zaitsev and Kislyakov, 2006). Acceleration in the super-Dreicer electric field driven by the R-T instability in the chromosphere foot points of the magnetic loops is more efficient (Zaitsev et al., 2016), but it requires high electric currents,  $I_0 > 10^{10}$  A, and the typical time of this process doesn't match the long-lived events. The model developed here can explain the electron acceleration in long-lived type IV radio continuum with the fine structures such as type III bursts, Zebra pattern, and sudden reductions and/or a bunches of type III bursts without type IV continuum background. Starting point of the type III bursts at 700 MHz mentioned by Huang et al. (2016) suggests that acceleration site is located near loop top and the beams of electrons directed toward the 3 GHz plasma frequency level. From Eq. 10 it follows also that pulsating type III events with period of about  $l/V_A \approx 1$  s (Aurass et al., 2003) can be explained in the lumped circuit approach if the current-carrying loop is quit compact,  $l \approx 10^8$  cm. Note that if  $v_{RLC}$  coincides with the frequency of the loop MHD oscillations, the ratio  $\Delta I/I_0$  grows and the acceleration and storage processes can be even more effective.

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