Astron. Astrophys. 320, 300–304 (1997)

ASTRONOMY
AND
ASTROPHYSICS

Magnetic reconnection mechanism for Type II white-light flares

X.Q. Li[1,2], **M.T. Song**[2], **F.M. Hu**[2], **and C. Fang**[3]

[1] CCAST (World Laboratory), P.O.Box 8730, Beijing, P.R. China
[2] Purple Mountain Observatory, Academia Sinica, Nanjing 210008, P.R. China
[3] Department of Astronomy, Nanjing University, Nanjing 210093, P.R. China

Received 17 July 1996 / Accepted 18 September 1996

Abstract. Based on the concept that the heating source for Type II white-light flares (WLFs) may be located in the photosphere, we suggest that Type II WLF is reduced by magnetic reconnection in a weakly ionized plasma. A plasmoid may rise by magnetic buoyancy from the subphotosphere into a sunspot region and then expand to form a current sheet in which a type of resistive instability takes place. The effect of neutral atoms on the resistive instability is estimated and found to produce a growth time of typically ten minutes that agrees well with observations.

Key words: magnetohydrodynamics (MHD) – plasmas – Sun: photosphere; flares

1. Introduction

In recent years, the existence of two types of white-light flare (briefly referred to as Type I and Type II) has been proposed by many authors (Boyer et al. 1985; Machado et al. 1986; Machado and Rust 1974; Mauas et al. 1990; Fang et al. 1994). For a Type I WLF, there is a strong Balmer jump in the spectra, and the Balmer lines are strong and very broad, and there exists a good time correlation between the maximum emission and peaks of hard X-ray and microwave radiation. On the contrary, for a Type II WLF, such a correlation does not exist, and the Balmer jump is very weak or disappears, and the Balmer lines are weak and relatively narrow. Through the computation of semi-empirical atmospheric models for the WLFs, it has been shown that the differences in the characteristic properties between Type I and II WLFs are due to the different temperature enhancements in the chromosphere and photosphere (Mauas 1990; Ryan et al 1983; Fang and Ding 1995): for the atmosphere of Type I WLFs there is a great enhancement of the chromospheric temperature but a small increase of the photospheric temperature, compared with the preflare atmosphere; on the contrary, for a Type II WLF the chromospheric temperature is not enhanced obviously while the photospheric temperature increases greatly. Therefore, they deduced that Type I WLFs are similar to normal major solar flares, that is, powerful electron beams bombard the chromosphere and then transport the energy to the photosphere; while for Type II WLFs the energy sources and the energy transport are quite different, and then it is possible that the heating source is located in the photosphere. Based on this idea, we suggest that a Type II WLF is due to magnetic reconnection phenomena in the photosphere which first leads to a temperature enhancement and then causes other effects accompanied with it.

After analysing the active regions associated with six Type II WLFs observed so far (Fang et al. 1994), it was found that three of them (the flares of 1982 June 15, 1972 Aug 7 and 1979 Sept 19) are located in usual isolated large δ-type sunspot groups, while the rest (the flares of 1980 July 1, 1984 May 20 and 1970 June 24) are within intermediate or small sunspot regions. But one thing is certain: all of them are located in magnetic turbulent regions with magnetic fluxes emerging frequently, and the magnetic configuration changing rapidly. Thus, it is reasonably assumed that Type II WLFs occur in magnetic turbulent regions, and that a plasmoid of about 4000 km in diameter with a magnetic field of 1000-2000 Gauss emerges from the subphotosphere (Such a plasmoid can be lifted out into a strong magnetic field region, especially near in the inversion line ($B_{\|} = 0$), and then it expands rapidly in the photosphere, forming a current sheet between the plasmoid and its surroundings (in more detail, see Hu et al. 1995). Within this current sheet, due to the coupling between the ponderomotive force and MHD, a new type of resistive instability takes place which leads to the release of a large amount of free magnetic energy. After the beginning of the reconnection process, another fast plasma instability ensues, repulsing the higher temperature particles out of the reconnection region, and then, as a result, the temperature remains about 5000 K (or rises slightly) while the majority of the magnetic energy is converted into that of excitation and ionization which produce the continuum radiation in the optical spectra, quickly leaving the source region.

In this paper we try to explain the following properties for a typical Type II WLF: the scale length of the flare kernel is $5 \sim 6$ arc sec $\approx 4 \cdot 10^3$ km, the lifetime is about 10 min. The

Send offprint requests to: X.Q. Li (second address above) densities of electrons and neutral atoms in the WLF kernel are 10^{13} and $10^{16}\mathrm{cm}^{-3}$, respectively. The power is 10^{28} erg·s^{-1} in optical wavelengths. It can be used to estimate the whole power: $10^{28} \cdot (1/0.9) \approx 1.2 \cdot 10^{28}\mathrm{erg}\cdot\mathrm{s}^{-1}$ which yields a total energy: $1.2 \cdot 10^{28} \times 600 \approx 7 \cdot 10^{30}$ erg. The temperature for a Type II WLF kernel is identified with that of upper photosphere, i.e. about 5000 K and the temperature enhancement is 150-250 K.

Since the degree of ionization in the photosphere is as low as 10^{-3} , it is necessary to consider the role of neutral gas in the reconnection process. In Sect. 2, under the assumption of a three-element (electrons, ions and neutral atoms) mixed fluid, the basic equations are first derived. And then in section a new type of linear resistive instability is investigated, which accounts for the main properties of Type II WLFs - growth rate, turbulent velocities, power etc. The computational results for a typical Type II WLF are presented in Sect. 4.

2. Derivation of basic equations

For a weakly ionized plasma, such as the photosphere, the friction force **R** caused by collisions between charged particles and neutral ones has to be considered. In such a case, the dynamics for electrons and ions are described by (Li et al. 1994; Zhang et al. 1995):

$$\frac{\partial n_e}{\partial t} + \nabla\cdot(n_e\mathbf{v}_e) = 0, \tag{1}$$

$$\frac{\partial n_i}{\partial t} + \nabla\cdot(n_i\mathbf{v}_i) = 0, \tag{2}$$

$$m_e n_e(\frac{\partial}{\partial t} + \mathbf{v}_e\cdot\nabla)\mathbf{v}_e$$

$$= e n_e(\mathbf{E} + \frac{1}{c}\mathbf{v}_e\times\mathbf{B}) - \nabla P_e + \mathbf{F}_p^e + R_{ei} + R_{en}, \tag{3}$$

$$m_i n_i\frac{\partial}{\partial t}\mathbf{v}_i + (\mathbf{v}_i\cdot\nabla)\mathbf{v}_i = -e n_i(\mathbf{E} + \frac{1}{c}\mathbf{v}_i\times\mathbf{B}) - \nabla P_i - R_{ei} + R_{in}, \tag{4}$$

where F_p^e is the ponderomotive force resulting from a high frequency field parallel to **B**:

$$\mathbf{F}_p^e = -m_e n_e\frac{1}{2}\nabla < (\mathbf{v}_f^e)^2 >, \tag{5}$$

and $\mathbf{v_f^e}$ is the fast oscillation velocity of electrons in the waves which is determined by the following transport equation:

$$\nabla\times\nabla\times\mathbf{v}_f^e + \frac{1}{c^2}\frac{\partial^2}{\partial t^2}\mathbf{v}_f^e + \frac{1}{c^2}\frac{4\pi e^2}{m_e}n_e\mathbf{v}_f^e = 0. \tag{6}$$

Since the collision terms associated with the neutral particles are related to their motions, the continuity and momentum equations for the neutral particles are:

$$\frac{\partial n_n}{\partial t} + \nabla\cdot(n_n\mathbf{v}_n) = 0, \tag{7}$$

$$m_n n_n\frac{\partial}{\partial t}\mathbf{v}_n + (\mathbf{v}_n\cdot\nabla)\mathbf{v}_n = -\nabla P_n - R_n. \tag{8}$$

For this three-element mixed fluid composed of electrons, ions and neutral particles, we can introduce a velocity related to the centre of mass:

$$\mathbf{u} = \frac{\rho_e\mathbf{v}_e + \rho_i\mathbf{v}_i + \rho_n\mathbf{v}_n}{\rho} \tag{9}$$

where ρ is the total mass density

$$\rho = \rho_e + \rho_i + \rho_n \equiv n_e m_e + n_i m_i + n_n m_n. \tag{10}$$

From Eqs.(1), (2) and (7), it follows that

$$\frac{\partial}{\partial t}\rho + \nabla\cdot(\rho\mathbf{u}) = 0. \tag{11}$$

Using the continuity equations for various particles, we can write down the following equality

$$n_\alpha m_\alpha(\frac{\partial}{\partial t} + \mathbf{v}_\alpha\cdot\nabla)\mathbf{v}_\alpha = \frac{\partial}{\partial t}(n_\alpha m_\alpha\mathbf{v}_\alpha) + \nabla\cdot(n_\alpha m_\alpha\mathbf{v}_\alpha\mathbf{v}_\alpha), \tag{12}$$

Then, adding Eqs.(2.3),(4) and (8) together, we obtain :

$$\rho(\frac{\partial}{\partial t} + \mathbf{u}\cdot\nabla)\mathbf{u} = \frac{1}{c}\mathbf{j}\times\mathbf{B} - \nabla P + \mathbf{F}_p^e, \tag{13}$$

where P is the total pressure related to the mass-centre system (Li 1990), and the current density **j** satifies the Maxwell equations:

$$\nabla\times\mathbf{B} = \frac{4\pi}{c}\mathbf{j}, \tag{14}$$

$$\nabla\times\mathbf{E} = -\frac{1}{c}\frac{\partial\mathbf{B}}{\partial t}, \tag{15}$$

$$\nabla\cdot\mathbf{B} = 0. \tag{16}$$

Multiplying (3) by (e/m_e), (4) by $(-e/m_i)$ and (8) by (e/m_n), and then adding them (also, using Eq.(12), we obtain the generalized Ohm law which can be simplified greatly when ignoring all terms associated with temporal and spatial derivatives and neglecting the Hall current as well as the small quantity (m_e/m_i), as has been done by Krall and Trivelpiece (1973). In this case, we have

$$\frac{n_s e^2}{m_e}\mathbf{E} + \frac{n_s e^2}{m_e c}(\mathbf{u}\times\mathbf{B}) + \mathbf{R} = 0, \tag{17}$$

where

$$\mathbf{R} = \frac{e}{m_e}\mathbf{R}_{ei}$$

$$+\frac{e}{m_e}\mathbf{R}_{en} - \frac{e}{m_i}\mathbf{R}_{in} - \frac{e}{m_n}\mathbf{R}_{in}, \quad (n_s \equiv n_e = n_i) \tag{18}$$

with

$$R_{\alpha\beta} = -m_\alpha n_s\nu_{\alpha\beta}(\mathbf{v}_\alpha - \mathbf{v}_\beta), \tag{19}$$

For the collision frequency between the charged particles and neutral ones, there is an approximation

$$\nu_{\alpha n} = n_n\sigma_{0\alpha}v_{T\alpha},$$

where $\sigma_{0\alpha}$ is the collision section-surface which is about $5 \cdot 10^{-15}$, depending weakly on the temperature, and therefore $\nu_{en} >> \nu_{in}$. Considering the average kinetic energy of charged particles satisfying $m_e v^2 \approx m_i v^2$ in thermal equilibrium, it follows that $v_e >> v_i$ or $\mathbf{j} \approx n_s e \mathbf{v}_e$. Assuming that the whole neutral particles form a stationary background, then $\mathbf{v}_n \approx 0$. Using the, one has

$$\mathbf{R} \approx -\nu_{ei}\mathbf{j} - \nu_{en} n_s e(\mathbf{v}_e - \mathbf{v}_n) \approx -(\nu_{ei} + \nu_{en})\mathbf{j}.$$

Hence, from Eq.(17), the generalized Ohm law can be recast in the form:

$$\mathbf{j} = \sigma(\mathbf{E} + \frac{1}{c}\mathbf{u} \times \mathbf{B}), \tag{20}$$

$$\sigma = \frac{n_s e^2}{m_e(\nu_{ei} + \nu_{en})}. \tag{21}$$

Combining Eqs.(11),(13),(5),(6), (20), (21),(14),(15) and (16) we obtain for the so-called three-element mixed fluid a set of equations coupling the ponderomotive force with the original MHD quantities:

$$\frac{\partial}{\partial t}\rho + \nabla \cdot (\rho \mathbf{u}) = 0, \tag{22}$$

$$\begin{aligned}\rho(\frac{\partial}{\partial t} + \mathbf{u} \cdot \nabla)\mathbf{u} \ &= \ \frac{1}{4\pi}(\nabla \times \mathbf{B}) \times \mathbf{B} - \nabla P \\ &\quad -\frac{1}{2}\frac{m_e}{m_i}\rho_i \nabla < (\mathbf{v}_f^e)^2 >, \end{aligned} \tag{23}$$

$$\frac{\partial}{\partial t}\mathbf{B} = \nabla \times (\mathbf{u} \times \mathbf{B}) + \frac{c^2}{4\pi}\eta \nabla^2 \mathbf{B}, \tag{24}$$

$$\nabla \times \nabla \times \mathbf{v}_f^e + \frac{1}{c^2}\frac{\partial^2 \mathbf{v}_f^e}{\partial t^2} + \frac{1}{c^2}\frac{4\pi e^2}{m_e} n_s \mathbf{v}_f^e = 0, \tag{25}$$

$$\nabla \times \mathbf{B} = \frac{4\pi}{c}\mathbf{j}, \tag{26}$$

$$\nabla \times \mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}, \tag{27}$$

$$\nabla \cdot \mathbf{B} = 0. \tag{28}$$

where ρ_i is the mass density of the electro-conductive fluid composed of electrons and ions; $\eta = \frac{1}{\sigma}$ is the electric resistivity.

3. Tearing mode theory

For convenience, we shall perform an incompressible analysis, following Furth et al. (1963). We assume that there is a transverse plasmon oscillating along the y-direction, and the unperturbed field is $\mathbf{B}_0$, perpendicular to the x-direction in a slab geometry. for the unperturbed state we assume:

$$\mathbf{u}_0 = 0, \quad \rho_0 = \bar{\rho} + \rho_0^1(x),$$

$$\mathbf{B}_0 = B_{0y}(x)\hat{y} = \bar{B}\ tanh(\frac{x}{L_s})\hat{y}.$$

Within the current sheet (i.e.$|\bar{x}| = |\frac{x}{L_s}| \ll 1$), the magnetic field can be approximated by

$$\mathbf{B}_0 = \bar{B}\frac{x}{L_s}\hat{y}, \tag{29}$$

where L_s is the characteristic length of the current sheet. From Eq.(23) the result is

$$\rho_0^1(x) \approx -\frac{1}{c_s^2}\frac{\bar{n}_s m_e}{2} < (\mathbf{v}_{f0}^e)^2 >, \tag{30}$$

when

$$\frac{1}{8\pi}\bar{B}^2(\frac{x}{L_s})^2 \ll \frac{\bar{\rho}_i}{2}\frac{m_e}{m_i} < (\mathbf{v}_{f0}^e)^2 >,$$

where c_s is the sound speed. Assuming that the high density neutral particles are considered as a non-mobile sea, one has $\rho_0^1(x) \approx \rho_{io}^1(x) \equiv m_i(\delta n_s)$. Then (30) becomes

$$\delta n_s = \frac{\rho_0^1(x)}{m_i} = -\frac{1}{c_s^2}\frac{\bar{n}_s \mu}{2} < (\mathbf{v}_{f0}^e)^2 >, \quad (\mu = \frac{m_e}{m_i}). \tag{31}$$

In this case Eq.(25) becomes

$$i\frac{\partial}{\partial t}\mathbf{v}_0 + \frac{c^2}{2\omega_{pe}}\nabla^2 \mathbf{v}_0 + \frac{\omega_{pe}}{8c_s^2}\frac{m_e}{m_i}|\mathbf{v}_0|^2 \mathbf{v}_0 = 0, \tag{32}$$

with

$$|\frac{1}{\omega_0}\frac{\partial}{\partial t} ln\ v_0| \ll 1.$$

where $\mathbf{v}_0 = v_0(x,t)\hat{y}$ is the amplitude of the fast oscillation:

$$\mathbf{v}_{f0}^e = \frac{1}{2}[\mathbf{v}_0(\mathbf{r},t)e^{i\omega_{pe}t} + c.c.],$$

with $\omega_{pe} \equiv \frac{4\pi \bar{n_s} e^2}{m_e}$, $n_s = \bar{n_s} + \delta n_s$, and c.c. representing the complex conjugae.

From Eq.(32) one finds that, due to the repulsion of the ponderomotive force, the density(30) within the current sheet is lower (Li and Wu 1989):

$$\rho_0^1(x) = -\frac{1}{4c_s^2}\bar{\rho}_i \mu (\bar{v}_0)^2 \mathrm{sech}^2(\frac{x}{\epsilon_0}), \tag{33}$$

where ϵ_0 is the width of the solitary wave:

$$\epsilon_0 = \frac{\sqrt{8}}{\sqrt{\mu}}(\frac{c}{\omega_{pe}})(\frac{c_s}{\bar{v}_0}).$$

We see from the above that the ponderomotive force by transverse plasma waves excludes plasma, resulting in a localized drop in density, as indicated by Eq.(33).

Examining the perturbation state in the form

$$A(x,y,t) = A(x)e^{ik_y y}e^{\gamma t},$$

and considering that $\rho \equiv \delta\rho \approx \delta\rho_i$, then from Eqs.(22) through (28), we obtain

$$\psi'' = \alpha^2 \psi(1 + \frac{\gamma \tau_R}{\alpha^2}) - ik_y \tau_R u_x F, \tag{34}$$

$$(\theta_0 u_x')' =$$

$$\alpha^2 u_x[\theta_0 + \frac{S^2}{\gamma^2\tau_R^2}G_0 + \frac{S^2F^2}{\gamma\tau_R}] + (\frac{i}{k_y\tau_R})\psi\alpha^2S^2(F - \frac{F''}{\gamma\tau_R}), \quad (35)$$

in which

$$B_{0y}(x) = \bar{B}F(x), \quad \psi = \frac{B_x}{\bar{B}},$$

$$\rho_0 = \bar{\rho}\theta_0(x), \quad x = L_s\bar{x}, \quad \alpha = kL_s, \quad S = \frac{\tau_R}{\tau_H}, \quad (36)$$

$$\tau_R = \frac{4\pi L_s^2}{\eta c^2}, \quad \tau_H = 4\pi L_s\frac{\sqrt{\bar{\rho}}}{\bar{B}}; \quad (37)$$

and

$$G_0 = \frac{\bar{\rho}_i}{\bar{\rho}}[\pi\mu^2(\frac{\bar{v}_0}{c_s})^2(\frac{\bar{v}_0}{\bar{v}_a})^2(\frac{L_s}{\epsilon_0})^4\bar{x}^2, \quad \bar{v}_A = \frac{\bar{B}}{\sqrt{4\pi\bar{\rho}}}, \quad (38)$$

where the superscript prime "$\prime$" represents the derivative with respect to $\bar{x}$. In the absence of transverse plasma waves, i.e. $G_0 = 0$, Eqs.(34) and (35) are reduced to those discussed by Furth et al.(1963).

Outside the current sheet ($\eta \to 0$), the solution for Eqs.(34) and (35) is (Li and Wu 1989)

$$\psi = \frac{i\ k_y\ F\ u_x}{\gamma}, \quad (39)$$

$$\psi_\pm = \begin{cases} e^{-\alpha\bar{x}}(1 + tanh\ \frac{\bar{x}}{\alpha}) & \text{if} \quad \bar{x} > 0, \\ e^{\alpha\bar{x}}(1 - tanh\ \frac{\bar{x}}{\alpha}) & \text{if} \quad \bar{x} < 0; \end{cases} \quad (40)$$

using the incompressible condition yields from (39) and (40)

$$u_y(x) = -\frac{\gamma}{k_y^2 L_s}\frac{1}{\bar{x}^2}, \quad (\bar{x} \ll 1);$$

then the average value over the current sheet scale L_s is

$$\bar{u}_y = \frac{1}{L_s}\int_{-\frac{L_s}{2}}^{\frac{L_s}{2}} u_y(\bar{x})\,d\bar{x} = \frac{4\gamma}{k_y^2 L_s}. \quad (41)$$

And one obtains the jump condition across the sheet in the outer region:

$$\Delta_e' \equiv \frac{(\psi_{+0}' - \psi_{-0}')}{\psi(0)} = 2(\alpha^{-1} - \alpha).$$

On the other hand, under a constant ψ_0- approximation, one can obtain the jump condition within the current sheet in a similar way given by Furth et al.(1963). By equating both jump conditions, one has (Li et al. 1994)

$$4\pi\frac{(\gamma\tau_R)^{\frac{5}{4}}\Gamma(\frac{3}{4})}{(\alpha s)^{\frac{1}{2}}\Gamma(\frac{1}{4})}(1+\xi)^{\frac{1}{4}} = 2(\frac{1}{\alpha} - \alpha). \quad (42)$$

with Γ being the gamma function and

$$\xi = \frac{d_0^2}{\gamma\tau_R}, \quad (43)$$

$$d_0^2 = \frac{\rho_i}{\bar{\rho}}[\pi\mu^2(\frac{\bar{v}_0}{c_s})^2(\frac{\bar{v}_0}{\bar{v}_A})^2(\frac{L_s}{\epsilon_0})^4], \quad (44)$$

Eq.(42) is the key equation for determining the resistive instability by solitary waves ($\xi >> 1$), and the growth rate of this instability can be deduced from Eq.(42) in the form

$$\gamma \approx \frac{1}{d_0^{\frac{1}{2}}}\frac{\alpha^{\frac{1}{2}}(\frac{1}{\alpha} - \alpha)s^{\frac{1}{2}}}{2\pi}\frac{\Gamma(\frac{1}{4})}{\Gamma(\frac{3}{4})}\frac{1}{\tau_R} \quad (\text{sec}^{-1}). \quad (45)$$

When the reconnction process occurs, the constant ψ_0- approximation is no longer valid; in this case a localized instability could be involved. By introducing the Fourier transformation as suggested by Furth et al. (1963), one obtains from Eqs.(34) and (35) (Li and Wu 1989; Li 1990)

$$\gamma_b \approx \frac{\bar{v}_A}{d_{eff}}, \quad (46)$$

with

$$d_{eff} = \frac{\sqrt{4\pi}}{d_0}\alpha L_s. \quad (47)$$

This instability determined by(46) is responsible for carrying away the products of field annihilation at a fast enough rate.

4. Discussion

For a typical Type II WLF (see Sect. 1), in its current sheet, the temperature,electron density and neutral atom density are $T = 5000K, n_e = 10^{13}\text{cm}^{-3}, \ n_n = 10^{16}\text{cm}^{-3}$ respectively; the magnetic field is taken to be $B = 10^3G$. These lead to the following plasma parameters: $v_{Te} = 3.89 \cdot 10^5T_e^{1/2} = 2.75 \cdot 10^7\text{cm}\cdot\text{s}^{-1}, \ \omega_{pe} = 5.64\cdot10^4n_e^{1/2} = 1.78\cdot10^{11}\text{cms}^{-1}, \ v_A = 2.18\cdot10^{11}n_n^{-1/2}B = 2.18\cdot10^6\text{cm}\cdot\text{s}^{-1}, \ c_s = \sqrt{3}\cdot3.89\cdot10^5\frac{T^1/2}{43} = 1.1\cdot10^6\text{cm}\cdot\text{s}^{-1}$. To estimate the amplitude $\bar{v}_0$ of fast electron oscillations $\mathbf{v}_{f0}^e$, it is convenient to assume that the wave-wave and wave-particle interactions lead to equipartition of energy over the both Langmuir and transverse plasmons. Then one has (Li and Ma 1994): $\bar{W}^p = (E_f^2/8\pi)/(n_ek_BT_e) > N_D^{-1} = \omega_{pe}^3/(n_ev_{Te}^3)$. Using the momentum equation, $\partial v/\partial t \approx em_e^{-1}\mathbf{E}_f$, yields $\bar{W}^p \approx (v_{f0}^e)^2/(2v_{Te}^2) = (\bar{v}_0^2)/(4v_{Te}^2) > N_D^{-1}$. So, with a large reserve, one may take $\bar{v}_0 = \sqrt(4v_{Te}^2N_D^{-1})$. It follows that the Debye number $N_D = 343\cdot T_e^{3/2}n_e^{-1/2} = 38.3, \ N_D^{-1} = 2.6\cdot10^{-2}, \ \bar{v}_0 = 9\cdot10^6\text{cms}^{-1}$. And then the width of the steady solitary wave is $\epsilon_0 = 2.5$cm.

The collision frequency between the electrons and ions is $\nu_{ei} \approx \omega_{pe}/N_D = 1.78\cdot10^{11}/38.3 = 4.65\cdot10^9s^{-1}$ and the collision frequency between electrons and neutrals is $\nu_{en} = n_n(5\cdot10^{-15})v_{Te} = 1.4\cdot10^9s^{-1}$. Then the conductivity is $\sigma = \bar{n}_se^2(m_e(\nu_{ei}+\nu_{en}))^{-1} = (\omega_{pe}^2/4\pi)\cdot10^{-9}/(6.05) = 4.2\cdot10^{11}s^{-1}$ and the resistivity $\eta = \sigma^{-1} = 2.4\cdot10^{-12}\ s$.

Computing the structure of current sheet's resistive instability: $\tau_R = 5.82\cdot10^{-9}L_s^2, \ \tau_H = 1.29\cdot10^{-6}L_s, \ S = \tau_R/\tau_H = 1.29\cdot10^{-6}L_s, \ d_0 = 7\cdot10^{-3}L_s^2, \ d_0^{-1/2} =$

$12L_s^{-1}$. The growth rate γ of the linear instability is computed as follows: $\alpha = L_s k_y = L_s 2\pi(\lambda_y)^{-1} << 1$, $\gamma = d_0^{-1/2}(2\pi)^{-3/2}L_s^{-1/2}S^{1/2}(3.63/1.23)\lambda_y^{1/2}\tau_R^{-1} = 1.74 \cdot 10^6 L_s^{-3}\lambda_y^{1/2}$, $\bar{u}_y = 1.76 \cdot 10^5 L_s^{-4}\lambda_y^{5/2}$. When taking $L_s = 1.5 \cdot 10^4$cm, $\lambda_y = 10^7$cm, we have $\gamma = 1.63 \cdot 10^{-3}$.

Then $\gamma^{-1} = 6.13 \cdot 10^2$sec ≈ 10 min, which is identified with the lifetime of a Type II WLF. It is noted that the y-direction motion shows alternative velocities(positive and negative) corresponding to the observational red and blue shift movement. The average velocity $(\bar{u}_y)_{max} = 1.1 \cdot 10^6$cms^{-1} is in good agreement with the observation.

The total energy released from the interaction region for a typical Type II WLF can be estimated as:

$$Q = \frac{c^2}{4\pi^2}\eta\frac{B^2}{d}\bar{A}\tau \approx \frac{5.4}{d} \cdot 10^{33},$$

where the current sheet area $\bar{A} = (4 \cdot 10^8)^2$ cm^2 and the lifetime $\tau = 613$ sec have been chosen. When the effective thickness of the sheet d is taken to be $7.7 \cdot 10^2$cm, we obtain $Q = 7 \cdot 10^{30}$ergs.

From Eq.(46), the growth rate of the fast plasma instability is $\gamma_b = \bar{v}_A/d_{eff} = 6.8 \cdot 10^9 s^{-1}$. Therefore the hot plasma, as a product of magnetic field dissipation, is quickly removed away in a very short period ($\gamma_b^{-1} \sim 1.5 \cdot 10^{-10} sec.$) and the temperature is kept about 5000 K or rises slightly only.

Compared with a fully ionized plasma, it is noted that the role of the neutral atom component in the reconnection process is: 1) diminishing the average Alfvén speed by 30 times; 2) decreasing the ratio S ($= \tau_R/\tau_H$) by 30 times; 3) increasing the resistivity by 10 times; 4) making a shorter wavelength oscillation (smaller by 5 times than one). These play important role in the resistive instability.

Acknowledgements. The authors are grateful to the Climbing Programme and the National Natural Science Foundation of China for partial supports.

References

Boyer R.,Machado M.E.,Rust D.M.,Sotirovski P., 1985, Solar Phys.98,255
Ding M.D.,Fang C.,Gan W.Q.,Okamoto T., 1994, ApJ 429, 890
Fang C.,Ding M.D., 1994, Progress in Astronomy(china), 12,100
Fang C.,Ding M.D., 1995, A&AS 110, 99
Fang C.,Ding M.D.,Okamoto T.,Gan W.Q., 1993, Sci.China (series A) 36,217
Furth H.P.,Killen J.,Rosenbluth M.N., 1963, Phys. Fluids 6, 459
Hu F.M.,Song M.T.,Li X.Q., 1995, Ap&SS 229, 325
Krall N.A.,Trivelpiece A.W., 1973, Principles of Plasma Physics, Mcgraw-hill Book Company, Sect. 3.2
Li X.Q.,Zhang Z.D.,Smartt R.N.,1994,A&A 290,963
Li X.Q.,Wu,S.T., 1989, in H.Kikuchi (ed.), Laboratory and Space plasma, Springer-verlag, New York, Inc., p.239
Li X.Q.,Ma,Y.H.,1994, Phys.Plasmas 1 (9), 1
Li X.Q., 1990, Publication of Purple Mountain Observatory, 9, 75
Machado M.E.,Avrett E.H.,Falciani R., 1986, in Neidig D.F.(ed.), Proc. of the National Solar Obs./Solar Max.Mission Symp., Sunspot, NSO, p.483
Machado M.E.,Rust D.M., 1974, Solar Phys. 38,499
Mauas P.J.,Machado M.E.,Avrett E.H.,1990, ApJ 360,715
Mauas P.J., 1990, ApJS 74, 609
Ryan J.M.,Chupp E.L.,Forrest D.J.et al., 1983, ApJ 272, L61
Zhang Z.D.,Li X.Q.,Smartt R.N., 1995, Ap&SS 226, 31

This article was processed by the author using Springer-Verlag LaTeX A&A style file *L-AA* version 3.