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Numerical analyses of self-generated magnetic fields with very small scale in solar corona

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Abstract. We numerically study the collapse behavior of self-generated magnetic fields described by nonlinear coupling equations. The results show that magnetic fields self-generated by transverse pumping plasmons may collapse, leading to an enhancement of the magnetic field of up to 250 – 500 Gauss in the range around 0.02 – 0.03 km, which corresponds to the inferred inhomogeneity small-scale fields in both the strength and scale in the solar corona.

Key words: Sun: corona – Sun: general – Sun: magnetic fields

1. Introduction

It is well known that there are cosmic magnetic fields in celestial bodies. The fields observed in galaxies, stars or in planets are large-scale macromagnetic fields. Besides the large-scale magnetic field in cosmic bodies, there are small-scale, non-uniform fields which cannot be resolved by present day instruments. Our sun is an example. Magnetic structures on the sun, excluding sunspots, are too small to explore directly (Stenflo 1989). It is shown that spatially highly intermittent flux fragments can occur all over the sun. There is now a considerable body of evidence suggesting that all scales of structure in the solar corona, as well as other objects of astrophysical interest, are coupled to small-scale processes associated with intermittent magnetic fields. Microwave spike bursts observed during solar flares exhibit characteristic properties of fine temporal structure, narrow bandwidth, high degrees of circular polarization and high brightness temperature (Slottje 1978; Benz 1986; Crannel et al. 1988; Gary 1991). It is believed that these properties are related to small-scale magnetic fields and their interactions. Mckean et al. (1990, 1989) investigated the source of radio spike bursts by using particle simulations to model the emission of maser radiation and its propagation through the solar corona. In their model, the formation of spike burst fine structure is dependent on the existence of small-scale magnetic field inhomogeneities in both the fundamental emission and the second-harmonic absorption layers. On the basis of the results of their simulation model, the source regions of the solar corona is located between $1264 < x/\triangle < 1864$ with which the non-uniform magnetic field has the characteristic scale of $l_c = (-1150{+}1984)\triangle = 834c/7\omega_{pe}$, where c is the velocity of light and ω_{pe} is the plasma frequency(from their Fig. 1), which corresponds to $l_c = 0.02\ km$ in coronal active regions(here $n_e = 10^9 cm^{-3}, \omega_{pe} = 5.64 \cdot 10^4\ n_e^{1/2}$). The non-uniform magnetic field is chosen so that ω_{pe}/Ω_e is about 0.4 – 0.2 (Ω_e is the electron gyrofrequency) or $B = \Omega_e/(1.76 \times 10^7) \approx 200\, Gs$ in coronal active regions. Obviously, however, the mechanism of convective instability (Parker 1978; Spruit 1979) is not at work.

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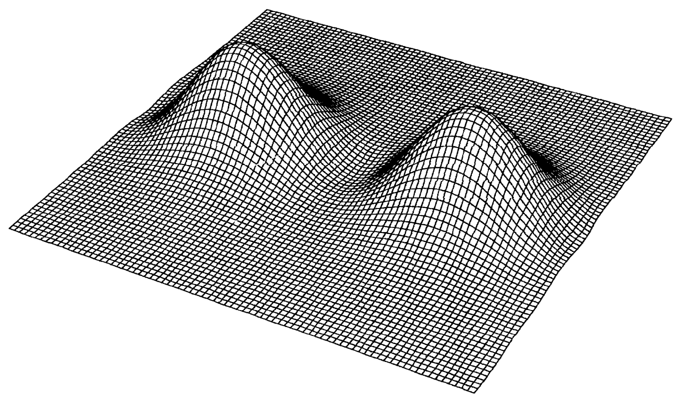

Fig. 1. The distribution of initial electric field

It is well known that magnetohydrodynamic (MHD) and two-fluid equations in solar corona follow from taking the first three moments of the Boltzmann equation, together with the macroscopic Maxwell equations; in the limits of large-scale and low-frequency motion the constitutive relations can be derived and the corona may be properly regarded as a well-behaved, multicomponent fluid. However, if the macroscopic length-scales of interest in the problem are not large and the dynamic frequencies in the problem are not low, the choice of appropriate constitutive relations and the concept of a local description of the equations of motion becomes questionable. At present we have studied the generation of localized magnetic flux with the characteristic scale $l_c < 1km$, which is very small compared to the electron-ion mean free path in the corona (Rosner et al. 1986):

$$l_{\rm mfp} \approx 1.1 \cdot 10^3 (\frac{n}{10^8})^{-1} (\frac{T}{10^6})^2 (\frac{\ln \Lambda}{20})\ {\rm km},$$

where $\ln \Lambda$ is the Coulomb logarithm, n is the density and T is the temperature; in such cases we may not define effectively the concept of a fluid element with the volume V and no rigorous perturbational expansion about the mean behavior will be possible under this circumstance. We must then generally attack the problem as one of kinetic theory rather than of fluid description.

Can one describe this very small-scale phenomenon by the use of two-fluid equations, including electron pressure? The answer is negative. For such cases of $l_c \ll l_{mfp}$, we could still formally introduce the concept of a fluid element, but the definition would be of little practical service because the particles that constitute the matter of the fluid element would freely stream into and out of the volume as a result of their random velocities. In other words, in these cases, a fluid description fails: we have no way of defining effectively the concept of a fluid element. Thus, it is not appropriate to study the generation of very small magnetic field by magnetohydrodynamic (MHD) and two-fluid equations.

On the other hand, we need to consider both high-frequency phenomena and interactions of plasma waves with particles in resonance, for example, the decay or coalescence processes of high-frequency fields and the modifications of particle densities by the low-frequency fields, as shown by Eqs. (14) and (19) in the next section. Hence, for the case of very localized regions during times of rapid changes, at the particle level, the kinetic theory would be required (Benz 1993): the fluid description for these processes is not valid.

Indeed, in solar corona, the number of particles within a Debye sphere, N_D is very large

$$N_D = \frac{4\pi}{3} n_e \lambda_e^3 \approx 1.4 \cdot 10^8 (\frac{n}{10^8})^{-0.5} (\frac{T}{10^6})^{1.5} \gg 1,$$

where $\lambda_e \equiv v_{Te}/\omega_{pe}$ is the Debye length; in other words the so called plasma approximation, required to make Debye shielding work and to justify a statistical treatment, is satisfied. Thus, the classical plasma kinetic theory is an appropriate description of the solar corona (Rosner et al. 1986). However, the processes involving rapid changes, small-scale and wave-particle interactions in resonance may be at the heart of some of the classical, unsolved coronal phenomena (Benz 1993). In addition, according to Mckean et al. (1989, 1990), the localized flux is attributed to a flare current system associated with the impulsive phase, i.e., to a neutral current sheet, where the background magnetic field is approximately zero with Maxwellian plasmas. And it has been shown analytically that the growth rate $\gamma_{\max}$ for linear instability of the self-generated magnetic field is much larger than the electron-ion collision frequency in hot and dilute corona (Li & Zhang 1996), implying that collisionless processes have turned out to control the situation. In summary, therefore, it is relevant to use the Vlasov equation and Maxwell equations, with vanishingly small background magnetic fields, to depict the wave-particle and wave-wave interactions of interest in the problem in the active region of corona in the next section. On the basis of the nonlinear equations derived by Li & Ma (1993), we try to examine numerically the self-generated magnetic fields with relevant initial values, resulting in the localized magnetic structures in coronal active regions in next sections.

2. Nonlinear equations for self-generated magnetic fields

We start from Vlasov's equation

$$\frac{\partial f_\alpha}{\partial t} + \mathbf{v}\frac{\partial f_\alpha}{\partial \mathbf{r}} + \mathbf{F}\frac{\partial f_\alpha}{\partial \mathbf{p}} = 0,\ (\alpha = e, i), \tag{1}$$

$$\mathbf{F} = e_\alpha [\mathbf{E} + \frac{\mathbf{v}}{c} \times \mathbf{B}], \tag{2}$$

where f_α is the particle distribution function

$$\int f_\alpha \frac{d\mathbf{p}}{(2\pi)^3} = n_\alpha, \tag{3}$$

where n_α is the particle density and $\mathbf{F}$ is the electromagnetic force. One can divide f_α and $\mathbf{F}$ into two parts:

$$f_\alpha = f_\alpha^R + f_\alpha^T, \qquad \mathbf{F} = \mathbf{F}^R + \mathbf{F}^T, \tag{4}$$

where $f_\alpha^R \left(\mathbf{F}^R\right)$ is unperturbed and $f_\alpha^T \left(\mathbf{F}^T\right)$ is perturbed. As f_α and $\mathbf{F}$ are closely coupled through Maxwell's equation and the current density equation

$$\mathbf{j} = \sum_\alpha \int e_\alpha \mathbf{v} f_\alpha \frac{d\mathbf{p}}{(2\pi)^3} \tag{5}$$

and the field $\mathbf{E}$ is assumed weak, so that the energy density excited is much smaller than the thermal one of the plasma, i.e.,

$$\bar{W} \equiv \frac{E^2}{8\pi n_e T_e} \ll 1, \tag{6}$$

we may expand f_α^T in powers of the perturbed field E^T

$$f_\alpha^T = \sum_\alpha f_\alpha^{T(a)}, \tag{7}$$

where the index a indicates that $f^{T(a)}$ is proportional to the a-th power of E^T. Put $\mathbf{F}^R = 0$ and expand $A = (\mathbf{F}, f)$ in a Fourier series

$$A(\mathbf{r}, t) = \int A_k e^{-i\omega t + i\mathbf{kr}} dk \tag{8}$$
$$(A_k = A_{\mathbf{k},\omega} \quad dk = d\mathbf{k} d\omega).$$

Taking Eqs. (5),(7) and the Maxwell equations into consideration, we can find the following field equation for the transverse mode from Eq. (1)

$$\left(k^2c^2 - \omega^2 \varepsilon_k^t\right) E_k^T = 4\pi i \omega \left[\mathbf{e}_\mathbf{k}^{t*} \cdot (\mathbf{j}_k^{(2)} + \mathbf{j}_k^{(3)})\right], \tag{9}$$

and the nonlinear currents (to the second and third order)

$$\mathbf{j}_k^{(2)} = \sum_\alpha \left(-e_\alpha^3\right) \int \frac{\mathbf{v} d\mathbf{p}}{(2\pi)^3} \frac{\hat{\mathbf{e}}_{\mathbf{k}_1}^t \cdot \frac{\partial}{\partial \mathbf{p}}}{\omega - \mathbf{k}\cdot\mathbf{v} + i\varepsilon} \cdot \left[\frac{\mathbf{e}_{\mathbf{k}_2}^t \cdot \frac{\partial}{\partial \mathbf{p}} f_\alpha}{\omega_2 - \mathbf{k}_2 \cdot \mathbf{v} + i\varepsilon}\right] \cdot E_{k_1}^T E_{k_2}^T \delta(k - k_1 - k_2) dk_1 dk_2, \tag{10}$$

$$\mathbf{j}_k^{(3)} = \sum_\alpha \left(+e_\alpha^4\right) \int \frac{\mathbf{v}d\mathbf{p}}{(2\pi)^3} \frac{i(\hat{\mathbf{e}}_{\mathbf{k}_1}^t \cdot \frac{\partial}{\partial \mathbf{p}})}{\omega - \mathbf{k}\cdot\mathbf{v} + i\varepsilon} \cdot \frac{\hat{\mathbf{e}}_{\mathbf{k}_2}^t \cdot \frac{\partial}{\partial \mathbf{p}}}{[(\omega-\omega_1) - (\mathbf{k}-\mathbf{k}_1)\cdot\mathbf{v} + i\varepsilon]} \cdot \left[\frac{\mathbf{e}_{\mathbf{k}_3}^t \cdot \frac{\partial}{\partial \mathbf{p}} f_\alpha^R}{\omega_3 - \mathbf{k}_3 \cdot \mathbf{v} + i\varepsilon}\right] \cdot E_{k_1}^T E_{k_2}^T E_{k_3}^T \cdot \delta(k-k_1-k_2-k_3) dk_1 dk_2 dk_3, \quad (11)$$

where ε_k^t is the transverse dielectric constant

$$\varepsilon_k^t = 1 + \sum_\alpha \frac{2\pi e_\alpha^2}{\omega k^2} \int \frac{(\mathbf{k}\times\mathbf{v})^2}{\omega - \mathbf{k}\cdot\mathbf{v} + i\varepsilon} \frac{\partial f_\alpha^R}{\partial \varepsilon_0} \frac{d\mathbf{p}}{(2\pi)^3}, \quad (12)$$

where ε_0 is the particle energy, $\mathbf{e}_\mathbf{k}^t$ the polarization vector for the transverse mode;

$$\hat{\mathbf{e}}_\mathbf{k}^t \equiv (1 - \mathbf{k}\cdot\mathbf{v}/\omega)\,\mathbf{e}_\mathbf{k}^t + \left(\mathbf{e}_\mathbf{k}^t \cdot \mathbf{v}\right)\mathbf{k}/\omega, \quad (13)$$

$$\delta(k - k_1 - k_2) \equiv \delta(\mathbf{k} - \mathbf{k}_1 - \mathbf{k}_2)\,\delta(\omega - \omega_1 - \omega_2), \quad (14)$$

the term $i\varepsilon$ in the denominators of integrated functions arises from the Landau rule. For the low-frequency fields $E_k^T = E_k^{Ts}$ and high frequency field $E_k^T = E_k^{T(+)}$, we obtain the following equations from Eq. (9):

$$\left(k^2c^2 - \omega^2\varepsilon_k^t\right) E_k^{Ts} = 4\pi i\omega \sum_\alpha \int \tilde{S}'^\alpha_{k,k_1,k_2} E_{k_1}^{T(+)} E_{k_2}^{T(-)} \cdot \delta(k-k_1-k_2)\, dk_1 dk_2, \quad (15)$$

$$\left(k^2c^2 - \omega^2\varepsilon_k^t\right) E_k^{T(+)} = 4\pi i\omega \sum_\alpha \{ \int \tilde{S}^\alpha_{k,k_1,k_2} E_{k_1}^{T(+)} \tilde{E}_{k_2}^{Ts} \cdot \delta(k-k_1-k_2)\, dk_1 dk_2 + \int \tilde{G}^\alpha_{k,k_1,k_2,k_3} E_{k_1}^{T(+)} E_{k_2}^{T(+)} E_{k_3}^{T(-)} \cdot \delta(k-k_1-k_2-k_3)\, dk_1 dk_2 dk_3 \}, \quad (16)$$

where $\tilde{S}^\alpha_{k,k_1,k_2}$, $\tilde{S}'^\alpha_{k,k_1,k_2}$ and $\tilde{G}^\alpha_{k,k_1,k_2,k_3}$ are matrix elements of interaction between high frequency fields and low frequency fields. The upper indicets "+" and "-" denote the positive and negative frequency parts for the high-frequency perturbation, and $\tilde{E}^{Ts}$ is different from E^{Ts} by a phase factor. Substituting Eq. (15) into Eq. (16) and retaining only the dominant contributions of electrons, we derive

$$i\frac{2}{\omega_{pe}} \frac{\partial \mathbf{E}(\mathbf{r},t)}{\partial t} - \frac{c^2}{\omega_{pe}^2} \nabla\times\nabla\times\mathbf{E}(\mathbf{r},t) = \frac{n'}{n_0}\mathbf{E}(\mathbf{r},t) + \frac{ie}{m_e c\omega_{pe}}\mathbf{E}(\mathbf{r},t)\times\mathbf{B}^*(\mathbf{r},t), \quad (17)$$

$$\left(\frac{\partial^2}{\partial t^2} - v_s^2\nabla^2\right) n'(\mathbf{r},t) = \nabla^2 \frac{|\,\mathbf{E}(\mathbf{r},t)\,|^2}{4\pi m_i}, \quad (18)$$

where $n'(\mathbf{r},t)$ is the perturbed density caused by the low-frequency fields

$$n_k'(\mathbf{r},t) = \left(n_k^{(1)} + n_k^{(2)}\right) = \frac{-k^2\left(\varepsilon_k^e - 1\right)}{4\pi m_e \omega_{pe}^2}\left[1 - \frac{\varepsilon_k^e - 1}{\varepsilon_k^l}\right] \cdot \int \mathbf{E}_{k_1}^{T(+)} \cdot \mathbf{E}_{k_2}^{T(-)} \delta(k-k_1-k_2)\, dk_1 dk_2, \quad (19)$$

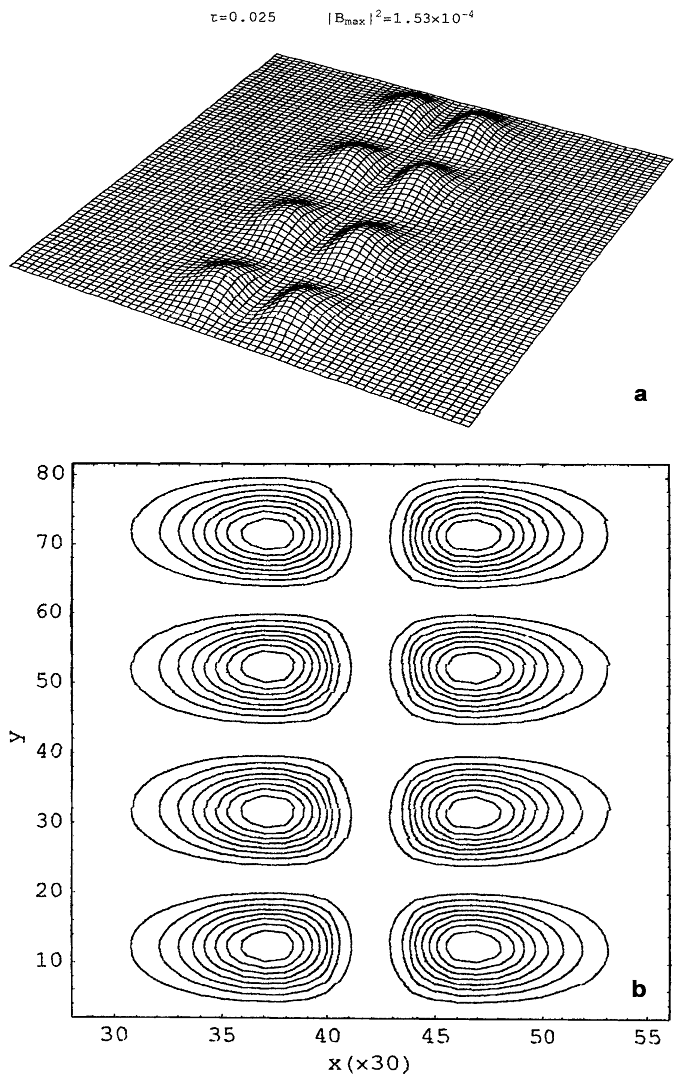

Fig. 2. Collapse development of self-generated magnetic field when $|E_{\max}|^2_{\tau=0} = 6$

and $\mathbf{E}(\mathbf{r},t)$ is the envelope of the high-frequency field

$$\mathbf{E}(\mathbf{r},t)\, e^{-i\omega_{pe}t} = \int \mathbf{E}_k^{T(+)} e^{-i\omega t + i\mathbf{k}\cdot\mathbf{r}} dk. \quad (20)$$

If the characteristic scale L of the low-frequency field is much larger than Debye length d_e ($L \gg d_e$), in coordinate representation, we can get the low-frequency field $\mathbf{B^s}$ from Eq. (15)

$$\left(-\frac{\partial^2}{\partial t^2} + v_s^2 \nabla\times\nabla\right)\mathbf{B^s}(\mathbf{r},\mathbf{t}) = \frac{iec}{m_e\omega_{pe}^2}\nabla\times\nabla\times\frac{\partial}{\partial t}\left[\left(\mathbf{E}(\mathbf{r},t)\times\mathbf{E}^*(\mathbf{r},t)\right)\right]. \quad (21)$$

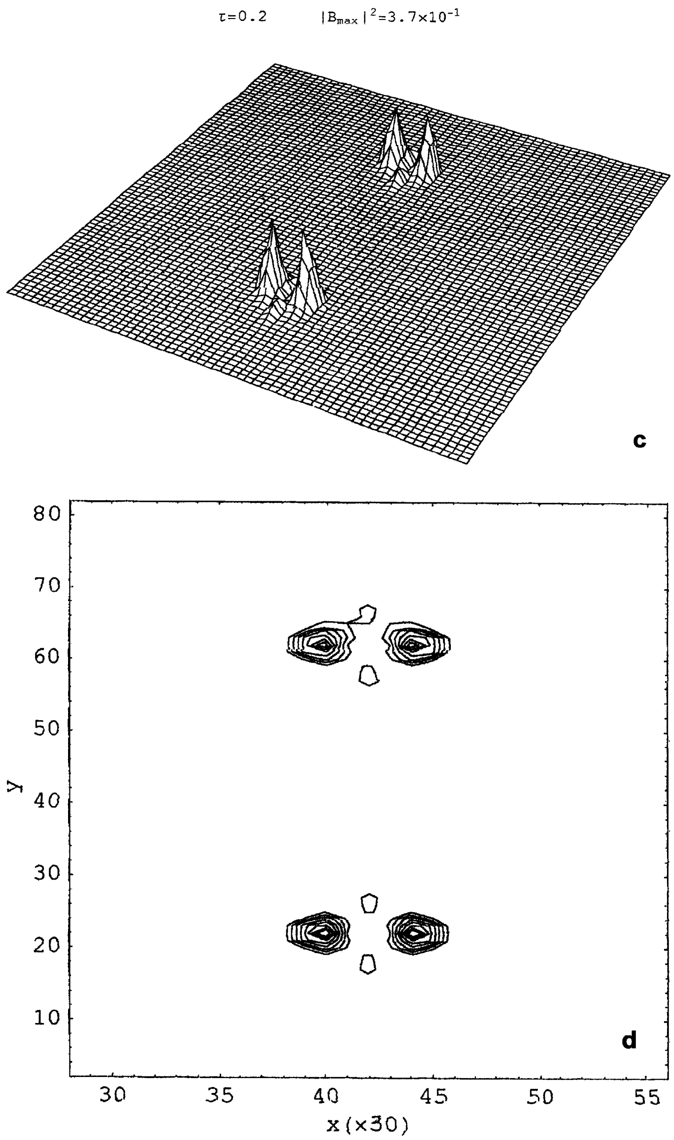

Fig. 2. (continued)

Through the substitutions

$$\xi=\frac{2}{3}\sqrt{\mu}\frac{1}{d_e}\mathbf{r},\quad \tau=\frac{2}{3}\mu\omega_{pe}t,\quad \mu=\frac{m_e}{m_i},\tag{22}$$

$$\alpha=\frac{c^2}{3v_{T_e}^2},\quad n=\frac{3}{4\mu}\frac{n'}{n_0},\tag{23}$$

$$\mathbf{E}(\xi,\tau)=\frac{\sqrt{3}\mathbf{E}(\mathbf{r},t)}{4(\pi\mu n_0T_e)^{1/2}},$$
$$\mathbf{B}(\xi,\tau)=\frac{3e}{4\mu m_e c\omega_{pe}}\mathbf{B}^S(\mathbf{r},t),\tag{24}$$

we can now write Eqs. (17),(18) and (21) in the following form (Li & Ma 1993)

$$\left(\frac{\partial^2}{\partial\tau^2}-\nabla^2\right)n(\xi,\tau)=\nabla^2\mid\mathbf{E}(\xi,\tau)\mid^2,\tag{25}$$

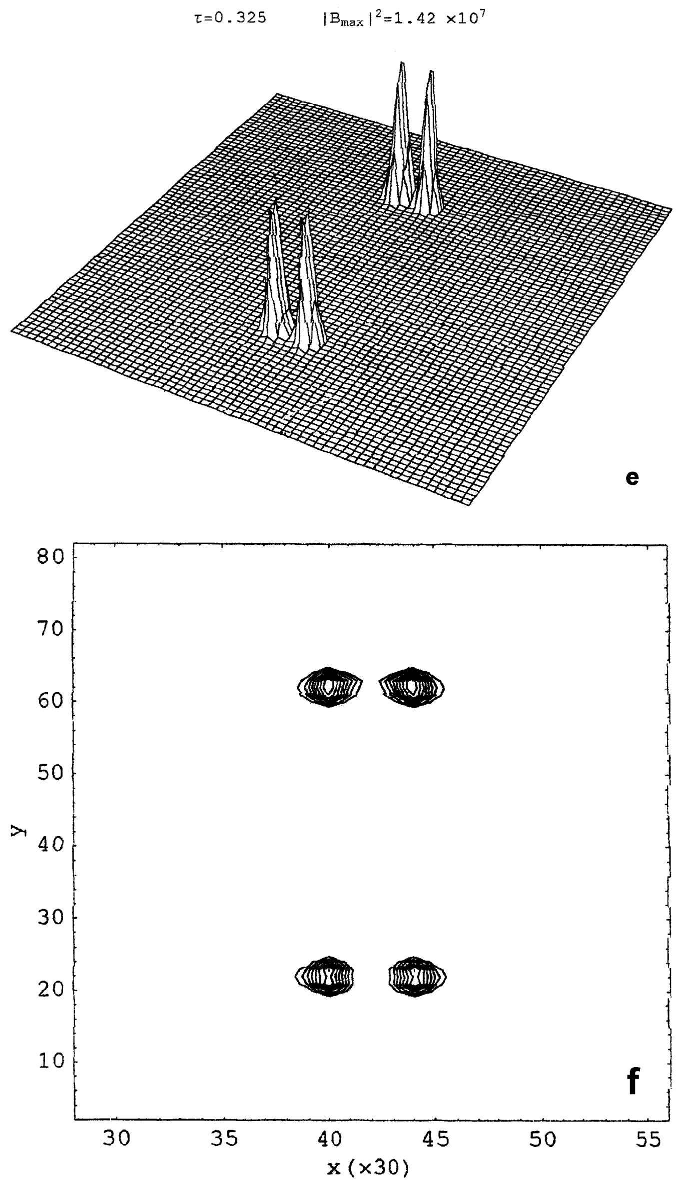

Fig. 2. (continued)

$$i\frac{\partial}{\partial\tau}\mathbf{E}(\xi,\tau)-\alpha\nabla\times\nabla\times\mathbf{E}(\xi,\tau)=$$
$$n(\xi,\tau)\mathbf{E}(\xi,\tau)+i\mathbf{E}(\xi,\tau)\times\mathbf{B}(\xi,\tau),\tag{26}$$

$$\left(-\frac{\partial^2}{\partial\tau^2}+\nabla\times\nabla\times\right)\mathbf{B}(\xi,\tau)$$
$$=i\frac{2}{3}\nabla\times\nabla\times\frac{\partial}{\partial\tau}\left[\mathbf{E}(\xi,\tau)\times\mathbf{E}^*(\xi,\tau)\right].\tag{27}$$

We restrict ourselves to the static limit for the low-frequency magnetic field and so Eq. (25) is reduced to

$$n=-|\mathbf{E}|^2.\tag{28}$$

Substituting Eq. (28) into Eq. (26) yields

$$i\frac{\partial}{\partial\tau}\mathbf{E}-\alpha\nabla\times\nabla\times\mathbf{E}+|\mathbf{E}|^2\mathbf{E}-i\mathbf{E}\times\mathbf{B}=0,\tag{29}$$

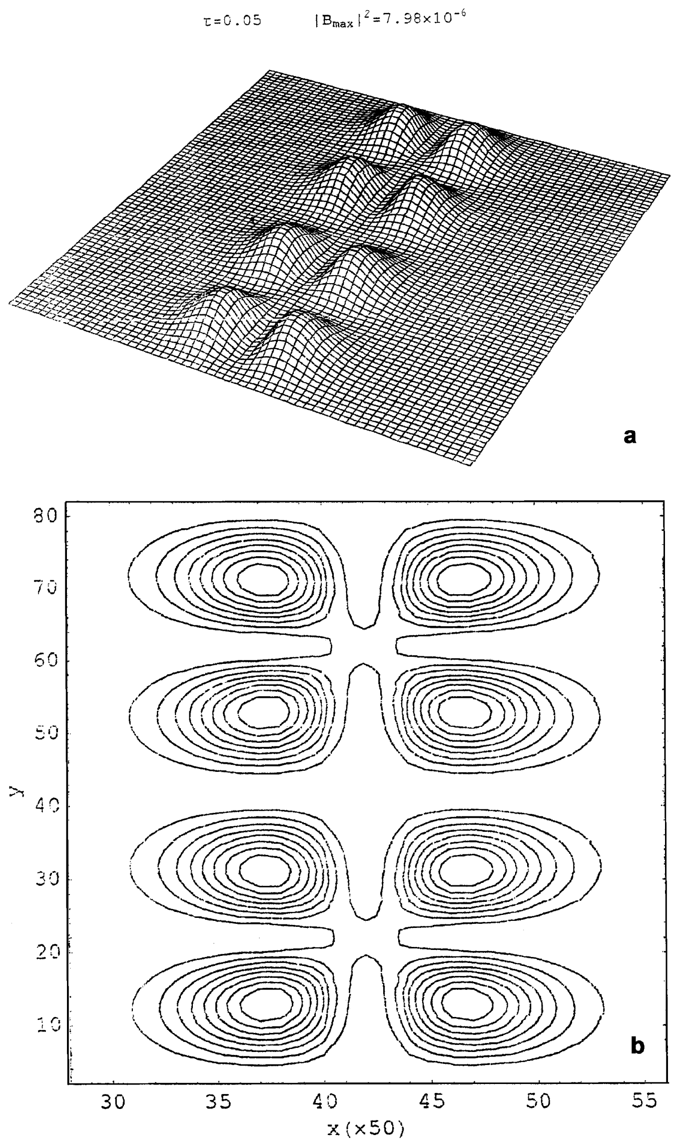

Fig. 3. Collapse development of self-generated magnetic field when$|E_{\max}|^2_{\tau=0} = 7$

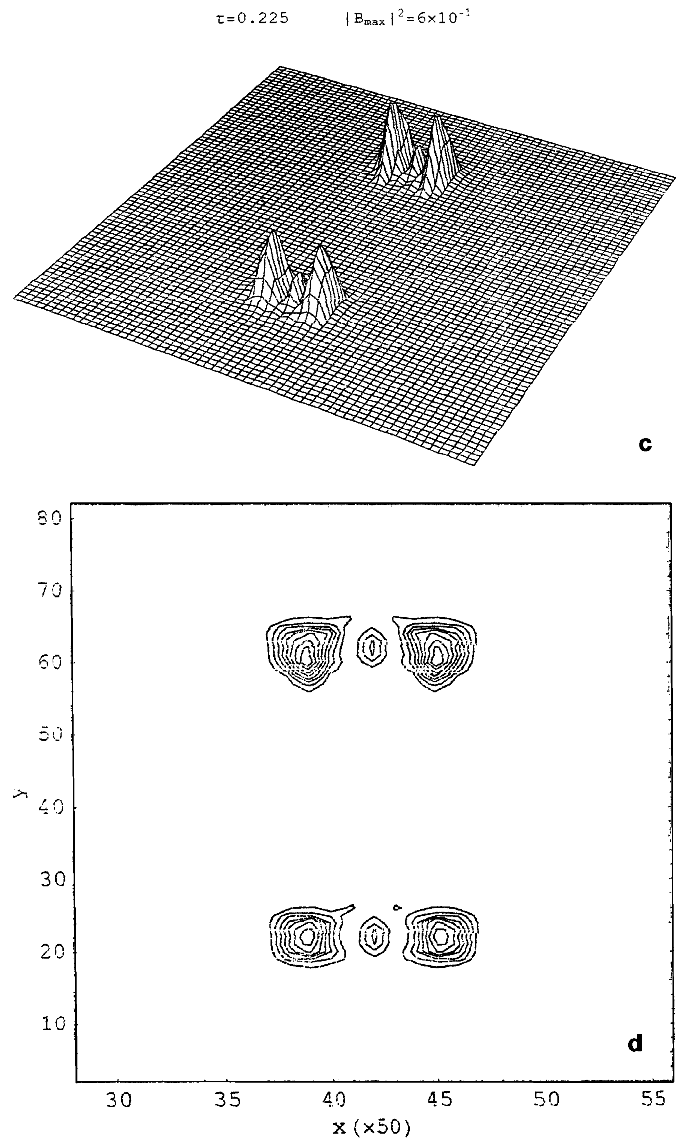

Fig. 3. (continued)

Eq. (27) is reduced to

$$\mathbf{B} = i\frac{2}{3}\frac{\partial}{\partial \tau}\left(\mathbf{E} \times \mathbf{E}^{*}\right). \tag{30}$$

Therefore it may be seen clearly that a self-generated magnetic field with very low-frequency is not an oscillatory field but a quickly collapsing field, which is completely determined by the closed Eqs. (29) and (30).

3. Numerical results

We have solved numerically Eq. (29) and Eq. (30) in two dimensions with three field components using a fast Fourier transform method. The problem with a periodic boundary-condition with respect to y is considered. The initial conditions which is satisfied to $\nabla \cdot \mathbf{E} = \mathbf{0}$ is chosen as

$$\mathbf{E}\left(\xi, \tau = \mathbf{0}\right) = \quad E_0 \sin(\frac{2\pi y}{y_0}) sech(\frac{x}{L_0})\mathbf{i} +$$

$$\left(-E_0 \frac{y_0}{\pi L_0^2}\right) \cos(\frac{2\pi y}{y_0}) th(\frac{x}{L_0}) sech(\frac{x}{L_0})\mathbf{j}$$
$$+E_0 \sin(\frac{2\pi y}{y_0}) sech\frac{x}{L_0}\mathbf{k}, \tag{31}$$

where y_0 is a period and L_0 is the width of the electromagnetic envelope. Evolution of the solution of Eq. (29) and Eq. (30) with the initial condition (31) is given in Figs. 2 – 4. These parameters are chosen as

$[E_{\max}]^2_{\tau=0} = 6,\ y_0 = 1 \times 10^3,\ L_0 = 3 \times 10^2$ (Fig. 2),

$[E_{\max}]^2_{\tau=0} = 6.12,\ y_0 = 1 \times 10^3,\ L_0 = 5 \times 10^2$ (Fig. 3),

$[E_{\max}]^2_{\tau=0} = 7.2,\ y_0 = 1 \times 10^3,\ L_0 = 5 \times 10^2$ (Fig. 4).

The distribution of the initial electric field is shown in Fig. 1. For coronal active regions, the numerical solution for Eqs. (29) and (30) show that the magnetic field self-generated by transverse

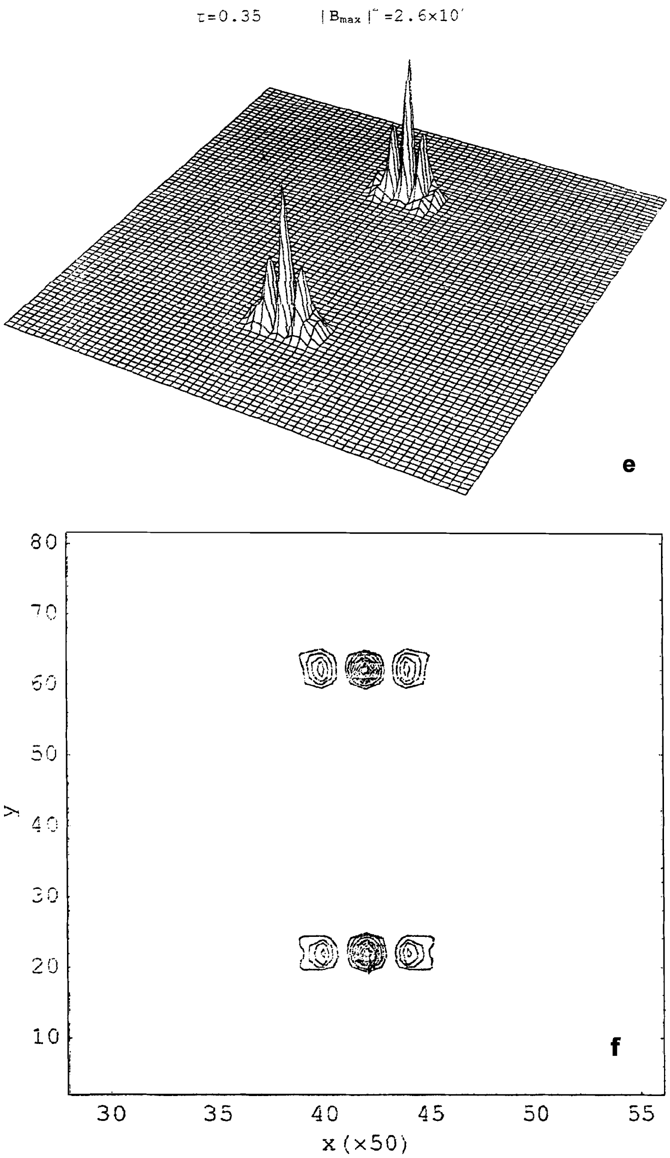

Fig. 3. (continued)

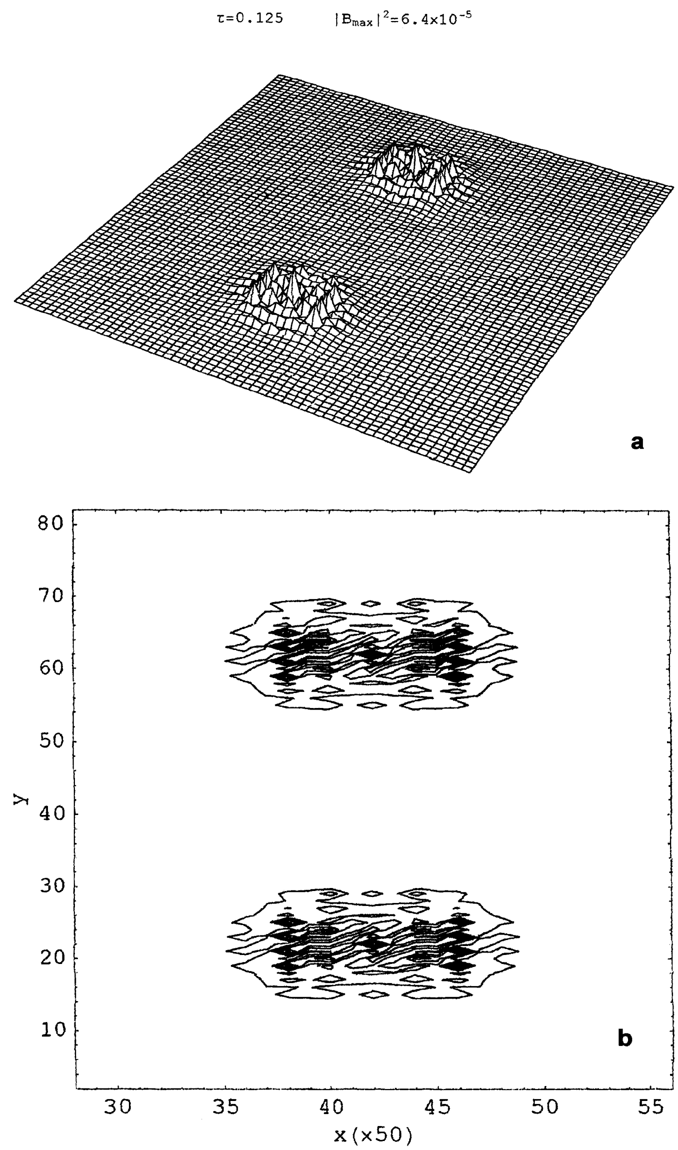

Fig. 4. Collapse development of self-generated magnetic field when$|E_{\max}|^2_{\tau=0} = 7.2$

plasmons with frequency $\omega \approx \omega_{pe}$ would collapse (see Figs. 2 – 4). The level contours of $|\mathbf{B}|^2 = B_x^2 + B_y^2 + B_z^2$ at successive times are shown in Figs. 2 – 4 for two dimensional geometry. In other words, due to self-compressing, a stronger magnetic field could be produced in a small region.Quantities in Figs. 1 – 4 are dimensionless. For coronal active regions, their relations to dimensional ones are(taking $n_e = 10^9$, $T_e = 3{\times}10^6$, $\alpha = 610$, $\omega_{pe} = 1.8 \times 10^9$)

$$B = 0.07(B)_{Fig}\,(G)\,,\ \frac{|E|^2}{4\pi n_e T_e} = 7.23 \times 10^{-4}\,|E|^2_{Fig}\,, \quad (32)$$

$$t = 1.6 \times 10^{-6}\tau\,(s)\,,\ \ l_c = 24.7\,(x)_{Fig}\,(\mathrm{cm})\,. \quad (33)$$

4. Conclusion and discussion

Low-frequency nonlinear plasma currents could be excited by high-frequency oscillation via wave-wave and the wave-particle interactions, leading to excitation of a very low-frequency magnetic field. It is shown that the dynamic behavior and configuration for the self-generated field is determined by Eqs. (29) and (30). If we take $|E_{\max}|^2_{\tau=0} = 6$ initially, then

$$\bar{W} \equiv \frac{E^2}{8\pi n_e T_e} = \frac{2m_e|E_{\max}|^2_{\tau=0}}{3m_i} = 2.1 \times 10^{-3}, \quad (34)$$

i.e., the turbulent electric field in dimension is $E = 4.4 \times 10^3$ V/m; the period and width are chosen as $y_0 = 1 \times 10^3$ and $L_0 = 3 \times 10^2$. Fig. 2 gives the resulting field strength and the characteristic size as $B^2_{\max} = 1.42 \times 10^7$ and $x = 1.08 \times 10^2$, respectively, and using Eqs. (32) and (33) yields

$$B = 263 \text{ Gauss}, \quad l_c = 2.7 \cdot 10^{-2} \text{ km}.$$

If we take $|E_{\max}|^2_{\tau=0} = 6.12$ initially, i.e., $E = 4.5 \times 10^3$ V/m, and take $y_0 = 1 \times 10^3$ and $L_0 = 5 \times 10^2$, the resulting

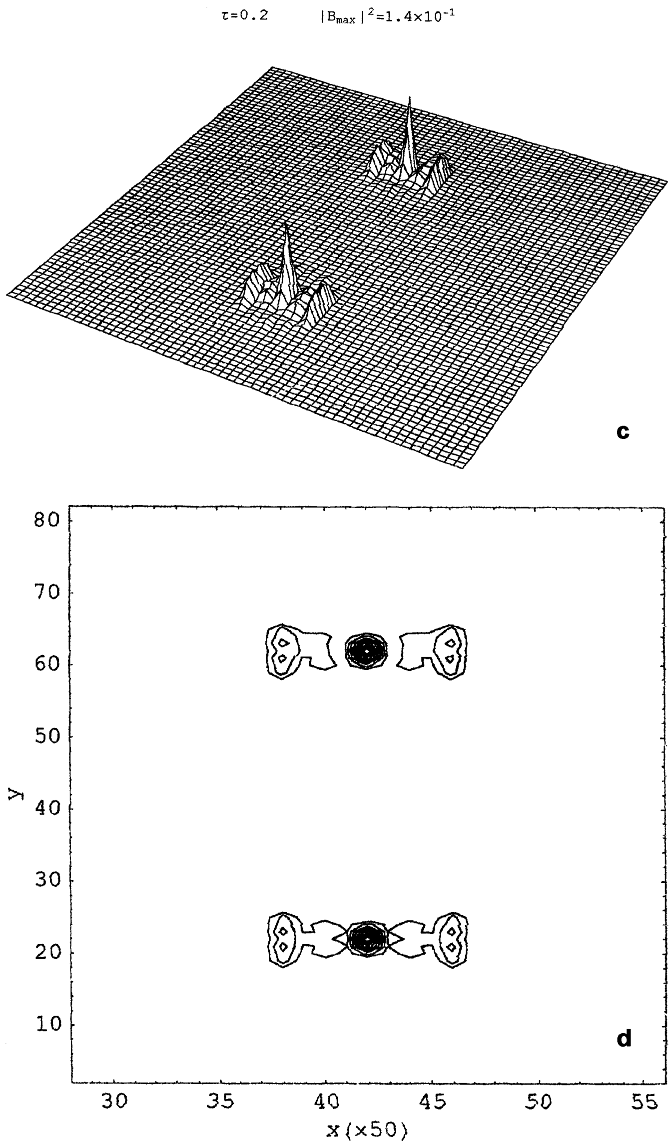

Fig. 4. (continued)

field strength and characteristic size are $B^2_{\max} = 2.6 \times 10^7$ and $x = 90$ (see Fig. 3), i.e.,

$$B = 360 \text{ Gauss}, \quad l_c = 2.2 \times 10^{-2} \text{ km}.$$

If we take $|E_{\max}|^2_{\tau=0} = 7.2$ initially, i.e., $E = 5.3 \times 10^3$ V/m, and take $y_0 = 1 \cdot 10^3$ and $L = 5 \times 10^2$, Fig. 4 gives the resulting field strength and characteristic size as $B^2_{\max} = 3.58 \times 10^7$ and $x = 87$, i.e.,

$$B = 420 \text{ Gauss}, \quad l_c = 2.1 \times 10^{-2} \text{ km}.$$

It is worth noting that Mckean et al. (1990, 1989) proposed a model for microwave spike bursts. In this model, the formation of fine structure is dependent on the existence of an intermittent magnetic field in coronal active regions. They inferred that the field would have a characteristic scale of 0.02 – 0.03 km (corresponding to $l_c = 834c/7\omega_{pe}$) and a strength of 250 – 500 Gauss (corresponding to $\omega_{pe}/\Omega_e = 0.4 - 0.2$), which are similar to

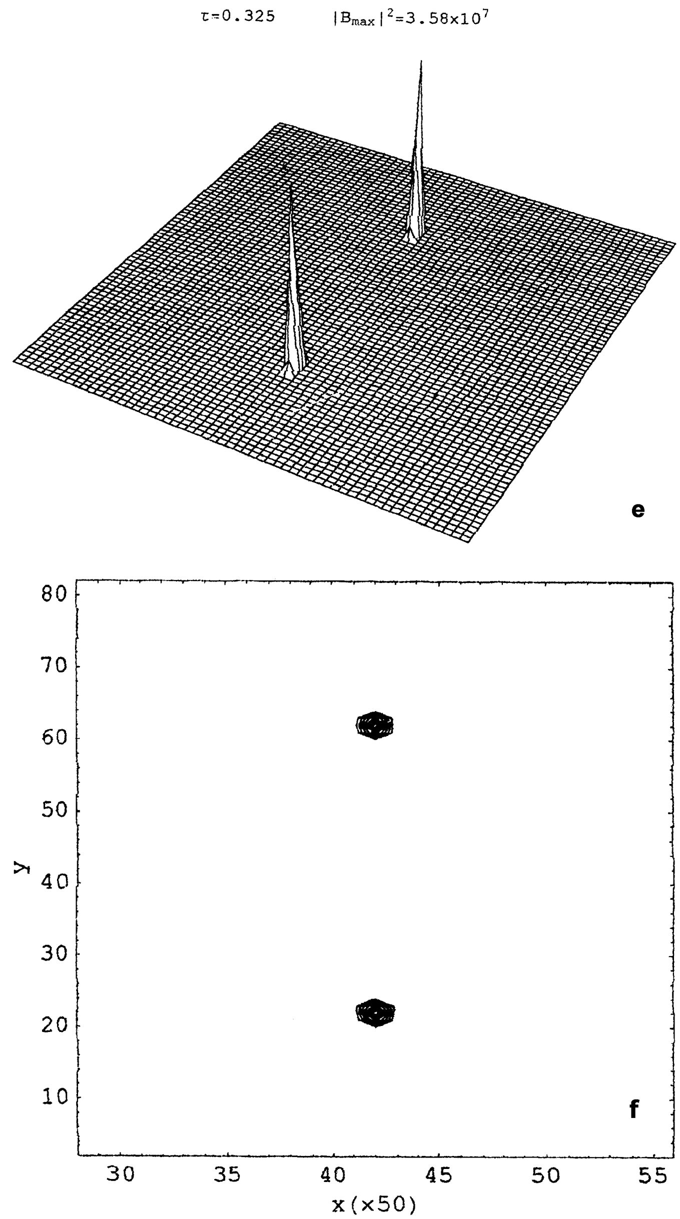

Fig. 4. (continued)

our results. In addition, it should be noted that the collapse tendency of self-generation magnetic fields is about the same for the three different initial values, although the resulting strength and characteristic size of such magnetic fields have slight difference in detail. This is a better result, which is not sensitive to the initial values.

On the other hand, the results of numerical analyses have shown that when $|E_{\max}|^2_{\tau=0}$ increases while L_0 remains unchanged, the collapse of self-magnetic fields become fast; when L_0 increases while $|E_{\max}|^2_{\tau=0}$ remains unchanged, the collapse of the self-magnetic fields becomes slow.

When the time scales are $\tau_1 > 0.325$ (see Fig. 2), $\tau_2 > 0.35$(see Fig. 3) and $\tau_3 > 0.325$(see Fig. 4) respectively, the field collapses rapidly and leads to a very strong field. $\bar{W} \equiv$

$E^2/(8\pi n_e T_e) = 2m_e|E_{\max}|^2/(3m_i) > 1$ and in this case the expansion Eq. (7) is no longer valid.

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