*ФИЗИКА ПЛАЗМЫ, 2016, том 42, № 5, с. 428–433*



*УДК 533.9*

# ONSET CONDITION OF THE SUBCRITICAL GEODESIC ACOUSTIC MODE INSTABILITY IN THE PRESENCE OF ENERGETIC-PARTICLE-DRIVEN GEODESIC ACOUSTIC MODE

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Received October 2, 2015

An analytic model is developed for understanding the abrupt onset of geodesic acoustic mode (GAM) in the presence of chirping energetic-particle-driven GAM (EGAM). This abrupt excitation phenomenon has been observed on LHD plasma. Threshold conditions for the onset of abrupt growth of GAM are derived, and the period doubling phenomenon is explained. The phase relation between the mother mode (EGAM) and the daughter mode (GAM) is also discussed. This result contributes to the understanding of 'trigger problems' of laboratory and nature plasmas.

**DOI:** 10.7868/S036729211605005X

## 1. INRODUCTION

The abrupt onset of large-amplitude deformation is often observed in laboratory and astro-space plasmas, such as disruptions in toroidal plasmas and burst of solar flares. One of the key issues is the mechanism, which gives a fast growth of the perturbation, like the fast reconnection in high temperature plasmas [1, 2]. The other essential problem is the sudden increase of the growth rate of the perturbation. For instance, the appearance of the $m = 1$ mode perturbation ($m$ is the poloidal mode number) at the beginning of the sawtooth crash has been reported in detail. For instance, the JET experiment [3] has reported that (1) the growth rate of the $m = 1$ mode perturbation is of the order of $(25\ \propto\mathrm{s})^{-1}$, and that (2) the growth rate increases to this large value within less than $100\ \propto\mathrm{s}$. That is, the growth rate itself increases with the growing rate which is of the order of the large growth rate. This abrupt onset of increase of the growth rate has been known as the 'trigger problem', and remains to be a mystery for more than a couple of decades.

The standard concept, which is constructed on the evolution of linear instabilities, is difficult to explain the trigger problem. The linear growth rate of a global mode is described by global equilibrium parameters, e.g. current profile, pressure gradient, etc., so that it changes with the time scale that governs global equilibrium parameters (which is denoted by $\tau$ in this article). When the linear stability boundary is crossed at $t = 0$, the linear growth rate evolves as $\gamma_L = \gamma_{L,0}(t/u) + \ldots$, where $\gamma_{L,0}$ is a characteristic value of growth rate in fully unstable cases. After the instability appears, the perturbation evolves as $X(t) \sim X(0)\exp(\gamma_{L,0}(t^2/2\tau))$. Therefore, the growth of the perturbation occurs with the time scale of $\sqrt{(\tau/\gamma_{L,0})}$. The relation $\sqrt{\tau/\gamma_{L,0}} \gg 1/\gamma_{L,0}$ usually holds, because $\tau$ belongs to global long time scales, so that the observation of abrupt increase of the growth rate (like [3]) cannot be explained. In the review [1], the working hypothesis based on the subcritical instability is explained to understand the trigger problem.

Recently, an abrupt onset of geodesic acoustic mode (GAM) [4, 5] has been observed on LHD [6, 7]. An energetic particle driven GAM (EGAM) [8–10] is observed in the NBI heated plasma in LHD [11]. The EGAM is associated with a frequency chirping, which is caused by the nonlinear instability mechanism associated with deformation of velocity distribution function [12]. It was found that when the frequency of the chirping EGAM becomes closer to the twice of the GAM frequency, the GAM is abruptly destabilized [7]. The growth rate is of the order of $(40\ \propto\mathrm{s})^{-1}$, and the growth rate jumps to this large value within a few $10\ \propto\mathrm{s}$. Thus, the understanding of this phenomenon will shed a light on the physics of trigger problem.

A physics model has been proposed [13], in which the subcritical instability of GAM in the presence of energetic particles [14] is combined with the parametric instability of EGAM–GAM system. The parametric instability induces the energy transfer from EGAM (mother mode) to GAM (daughter mode) so as to make a seed of GAM. If this seed amplitude is larger than the threshold of subcritical instability, nonlinear growth rate gives a large growth rate without modifying the global plasma parameters. Direct nonlinear simulation has been performed. The comparison with the experiments seems encouraging. The sensitivity of the simulation result to the parameters has been analyzed. In this article, we discuss the analytic model for the abrupt onset of GAM in the presence of chirping EGAM. Threshold condition is derived, and the period doubling phenomenon is explained. The phase relation between the mother and daughter modes are also discussed.

## 2. MODEL

We first summarize the experimental observation on LHD about the sudden onset of GAM, in order to clarify the motivation to construct the model. Next, the basic model equation is revisited, which was formulated as nonlinear kinetic equation in [13, 15]. Then the linearized model is deduced, in order to study the condition of the onset of daughter mode in the presence of the chirping mother mode.

### *2.1. Short summary of experimental observation*

The experimental observations, which form the basis of the model, are described in [7, 11]. The LHD is a helical system (major ($R$) and averaged minor radii of plasma are 3.75 m and about 0.6 m, respectively). In this experiment, the magnetic field is 1.375 T at the magnetic axis, and the ion species is hydrogen. The plasmas are sustained by a neutral beam injection (NBI) (the injected beam energy is 175 keV, and the absorbed power is about 140 kW). The line averaged electron density is $0.1\times10^{19}$ m$^{-3}$, the central electron temperature (heated by the electron cyclotron heating with a power of 2.5 MW) is approximately 8 keV, the ion temperature is several 100 eV, and the slowing-down time of the injected ion beam is about 20 s. A positive gradient in the velocity space of the beam is realized for strong EGAM excitation.

Under this circumstance, the EGAMs (frequency of which is in the range of a few tens kHz to 100 kHz) are observed. The identification has been done carefully. The potential and density perturbations were measured by ion beam probe, and the correlations with the signals from arrays of magnetic probes are evaluated. The toroidal mode number is found to be zero. The poloidal mode number is zero for potential perturbation and one for density perturbation. The frequency was compared with the theory with plausible distribution function of fast ions. With these examinations, the excitation of EGAM was identified [11]. When the mode is excited, the chirping of frequency to higher frequency is often observed. The time scale of one burst of chirping EGAM is of the order of 10 ms (which we refer to 'slow' in this article). The envelope grows and decays in this time scale as the frequency chirping takes place.

The most striking discovery is that during a slow burst of chirping EGAMs, a very sharp burst (time scale of which is a few 100 ∝s) happens. This time scale is referred to 'fast' here. If it occurs, the amplitude of oscillation becomes a few times larger, and as a result of this large perturbation, the change in the ion measurement (like effective temperature) is observed [16]. Such a sharp burst has a frequency close to the GAM eigenmode. The onset occurs, when the frequency of the chirping EGAM is close to twice of that of excited GAM. In addition, when the GAM appears, the phase relation between the EGAM and GAM is not random but preserved, so that the onset appears as a period doubling bifurcation of EGAM. In addition, the growing rate of the growth rate of GAM is close to the peak growth rate of GAM.

These observations stimulate to construct a model as follows: Since the GAM appears abruptly (and the mean plasma parameters, that determine the linear growth rate, change only very little during a short time of onset), we consider that the GAM is not driven by linear instability mechanism. The intensity of microscopic turbulence, which may drive GAMs, does not change substantially during this short time. Thus, we choose the hypothesis that the observed GAM may be induced by subcritical mechanism, which has been theoretically pointed out. For the subcritical excitation, some seed is necessary [17]. Considering the observation that the excitation takes place when the frequency of the chirping EGAM is close to twice of the excited GAM frequency, thepossibility of transferring the energy from EGAM to GAM via parametric coupling is investigated in the model.

### *2.2. Brief description of nonlinear model*

Here we consider the interaction of two modes. In order to treat the present problem, we split the electric field $E$ between the two waves, $E = E_1 + E_2$, and introduce a hybrid model. The daughter mode ($E_1$) is treated by the kinetic 1D model, and the mother mode ($E_2$) is treated as a simple medium for nonlinear energy transfer. For $E_2$, we prescribe the initial amplitude $Z_{2,0}$ and the time evolution of frequency $\omega_2(t)$ from the experimental data. We assume that the impact of the mother on particles near the resonant interaction of the daughter is negligible. The interaction between the two modes is modeled by the equation of period doubling. The evolution of the energetic particle distribution $f(x,v,t)$, in the neighborhood of the resonance of the daughter mode $E_1$, is given by a kinetic equation [18, 19]

$$\frac{\partial f}{\partial t}+v\frac{\partial f}{\partial x}+\frac{qE_1}{m}\frac{\partial f}{\partial v}=\frac{\nu_f^2}{k_1}\frac{\partial \delta f}{\partial v}+\frac{\nu_d^3}{k_1^2}\frac{\partial^2 \delta f}{\partial v^2}, \quad (1)$$

where $\delta f \equiv f - f_0$, and $f_0(v)$ is the initial velocity distribution. The right hand side is a collision operator, where $\nu_f$ and $\nu_d$ are input parameters characterizing dynamical friction and velocity-space diffusion, respectively. Thisoperator is obtained by projecting a Fokker–Planck operator that describes Coulomb collisions perceived by energetic ions, on the resonant phase space surface [19].

The evolution of the two parts of electric field is given by

$$\frac{dZ_1}{dt}=-\frac{m\omega_p^3}{4\pi q n_0}\int f(x,v,t)e^{-i(k_1x-\omega_1 t)}dxdv - \gamma_{damp}Z_1 - i\frac{V}{\omega_1}Z_2Z_1^*e^{i\theta t}, \quad (2)$$

$$\frac{dZ_2}{dt}=-i\frac{V}{\omega_2}Z_1^2e^{-i\theta t}, \quad (3)$$

where

$$E_j \equiv Z_j \exp[i(k_jx-\omega_jt)] + c.c., \quad (4)$$

$n_0$ is the total density, and $\theta = 2\omega_1 - \omega_2$ is the mismatch frequency. The term proportional to $\gamma_{damp}$ is an external wave damping, which is a model for all linear dissipative mechanism of the wave energy to the background plasma [18]. It is assumed to be constant in time in this model.

### 2.3. Linearized model

The subject of the article is to study the onset condition of the subcritical GAM instability in the presence of energetic particle driven GAM which is chirping in frequency. The parametric instability of the daughter mode in the presence of mother mode, the frequency of which is close to the twice of the daughter mode, is investigated. For the study of parametric instability of the daughter mode, we take the amplitude of the daughter mode as a smallness parameter. When the daughter mode is excited by parametric instability, the amplitude is much below the subcritical nonlinear destabilization, and the kinetic process (the first term in the RHS of Eq. (2)) is represented by its linear limit as

$$-\frac{m\omega_p^3}{4\pi q n_0}\int f(x,v,t)e^{-i\zeta_1}dxdv = \gamma_{L,0}Z_1, \quad (5)$$

where $\zeta_1 = k_1x - \omega_1 t$. Combining this kinetic growth process and other damping process, the linear damping rate of the daughter mode is represented by one parameter as

$$-\gamma_d = \gamma_{L,0} - \gamma_{damp}. \quad (6)$$

The amplitude of the mother mode is the zeroth-order term, and the correction due to the growth of the daughter mode is the second order correction. The frequency mismatch is slowly varying in time, so that it is assumed to be constant in the process of the growth of the daughter mode. With these consideration of the ordering, we study the time evolution of $Z_1(t)$ and $Z_1^*(t)$, and other terms like $Z_2$ and $\theta$ are treated as constant. Based on these analytical assumptions, one has a linearized equation as

$$\dot{Z}_1 = -\gamma_d Z_1 - \frac{iV}{\omega_1}Z_2Z_1^*e^{i\theta t}. \quad (7)$$

## 3. CONDITION FOR PARAMETRIC EXCITATION

### 3.1. Eigenvalue equation

Based on Eq. (7), the onset of parametric instability of the daughter mode is analyzed. Taking the time derivative of Eq. (7), one has

$$\ddot{Z}_1 = -\gamma_d\dot{Z}_1 + \frac{V\theta}{\omega_1}Z_2Z_1^*e^{i\theta t} - \frac{iV}{\omega_1}Z_2\dot{Z}_1^*e^{i\theta t}. \quad (8)$$

Time derivative $dZ_1^*/dt$ is given by taking the complex conjugate of Eq. (7), and is substituted into Eq. (8) to give

$$\ddot{Z}_1 = (-2\gamma_d + i\theta)\dot{Z}_1 + (i\theta\gamma_d - \gamma_d^2 + D^2)Z_1, \quad (9)$$

where $D^2$ is defined as

$$D^2 = \frac{|V|^2}{\omega_1^2}|Z_2|^2. \quad (10)$$

By imposing the time dependence

$$Z_1 \propto e^{\lambda t} \quad (11)$$

in Eq. (9), one has the eigenvalue equation

$$\lambda^2 + (2\gamma_d - i\theta)\lambda + \gamma_d^2 - i\theta\gamma_d - D^2 = 0. \quad (12)$$

The eigenvalue is obtained as

$$\lambda = -\gamma_d + \frac{i\theta}{2} \pm \sqrt{D^2 - \frac{\theta^2}{4}}. \quad (13)$$

### 3.2. Instability condition

Equation (13) gives the condition for the onset of parametric instability as

$$\frac{\theta^2}{4} < D^2 - \gamma_d^2. \quad (14)$$

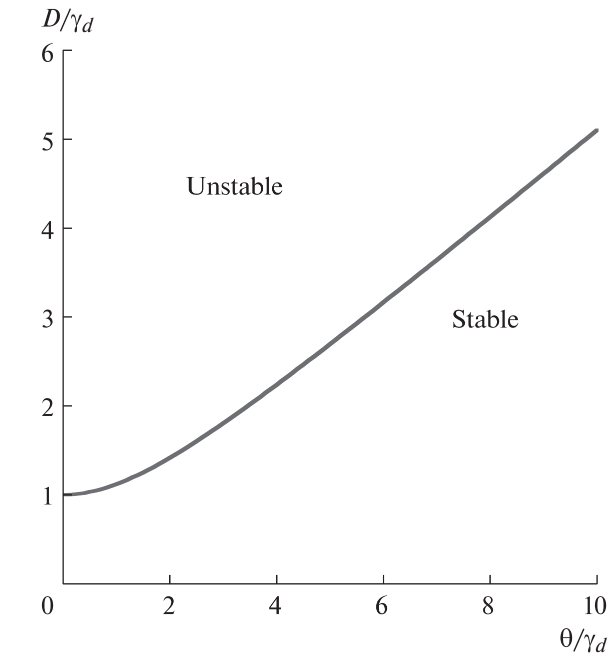


**Fig. 1.** Condition for the onset of parametric instability.

Figure 1 illustrates the domain of parametric instability of the daughter mode. The necessary condition for the parametric instability is that

$$D > \gamma_d. \tag{15}$$

The mother mode has a threshold, above which the parametric instability is induced. Once the mother mode exceeds the necessary amplitude so that Eq. (15) is satisfied, the parametric instability takes place when the frequency mismatch is small as to satisfy Eq. (14). It explains the observations in experiment and simulation that the onset of daughter mode takes place when the frequency of the mother mode is slightly lower than the twice of the daughter mode frequency. The boundary of domain of excitation of the daughter mode in Fig. 3 of [13] is close to $D \sim \gamma_d$ in the limit of large $D$, confirming Eq. (14). The onset of excitation of the daughter mode occurs even if the frequency mismatch does not vanish.

### *3.3. Phase-locking and period doubling*

The eigenvalue Eq. (13) gives

$$Im\lambda = \frac{i\theta}{2}. \tag{16}$$

That is, $Z_1$ is oscillating at the half frequency of the mismatch, and oscillatory part is given as

$$Z_1 \propto \exp\left(\frac{i\theta}{2}t\right). \tag{17}$$

The oscillatory part of the daughter mode has a half frequency of the mother mode as

$$\begin{aligned} E_1 &= Z_1 e^{-i\omega_1 t + ik_1 x} + c.c. \propto \\ &\propto \exp\left(-\frac{i\omega_2}{2}t + ik_1 x\right) + c.c., \\ E_2 &\propto \exp(-i\omega_2 t + ik_2 x) + c.c. \end{aligned} \tag{18}$$

For any value of the frequency mismatch $\theta$, the daughter mode has a half frequency of the mother mode. Thus the period doubling of the mother mode happens when the daughter mode is excited via parametric instability.

The phase of the daughter mode (with respect to the mother mode) is also obtained. We take the origin of time so that the variable $Z_2$ is real. Here we write

$$\frac{iV}{\omega_1}Z_2 = iD_0 e^{i\beta} \tag{19}$$

where $D_0 = \sqrt{D^2}$ and $\beta$ indicates the phase of the coupling constant $V$. The phase difference $\beta$ is considered to be small [20]. With this choice, and substituting Eqs. (11) and Eq. (13) into Eq. (7), one has

$$\left(\sqrt{D^2 - \frac{\theta^2}{4}} + \frac{i\theta}{2}\right)Z_1 = D_0 e^{-i\frac{\pi}{2}+i\beta} Z_1^* e^{i\theta t}. \tag{20a}$$

Let us abbreviate as

$$\frac{\sqrt{D^2 - \theta^2/4} + i\theta/2}{D_0} \equiv e^{i\alpha} \tag{20b}$$

to rewrite Eq. (20a) as

$$Z_1 = e^{(-\pi/2-\alpha+\beta)i} Z_1^* e^{i\theta t}. \tag{20c}$$

Introducing $Z_0$ to indicate the phase

$$Z_1 = Z_0 e^{\lambda t} \tag{21a}$$

in Eq. (20c), one sees that the daughter mode has a relative phase with respect to $Z_2$ as

$$Z_0 = |Z_0| e^{-i\delta} \tag{21b}$$

with

$$\delta = \frac{\pi}{4} + \frac{\alpha}{2} - \frac{\beta}{2}. \tag{21c}$$

The daughter mode has a relative phase of $\delta$, and is in advance with respect to the mother mode. In the case of marginal condition of $D \sim \theta/2$, one has $\alpha \sim \pi/2$. If $D \gg \theta$ holds, $\alpha \sim 0$. The phase difference $\delta$ is between $\pi/2$ and $\pi/4$.

### *3.4. Back-interaction on the mother mode*

One can also study the quasilinear response of the mother mode after the onset of daughter mode. From Eq. (3), one has

$$Z_2^* \frac{dZ_2}{dt} = -\frac{iV}{\omega_2} Z_2^* Z_1^2 e^{-i\theta t}. \tag{22}$$

The complex conjugate of Eq. (22) is added to Eq. (22) to have

$$\frac{d}{dt}|Z_2|^2 = -2Im\left(\frac{V^*}{\omega_2}Z_2Z_1^{*2}e^{i\theta t}\right). \tag{23}$$

The amplitude of the mother mode is subject to the second order correction owing to the growth of the daughter mode. The LHS is the second order term (with respect to the daughter mode), so that the $Z_2$ term in the RHS is evaluated by the zero-th order term, which is the original amplitude. We have

$$\frac{d}{dt}|Z_2|^2 = -Re\Gamma|Z_1|^2 = -\frac{1}{2}\frac{d}{dt}|Z_1|^2. \tag{24}$$

Here $Re\Gamma \equiv 2Im((V^*/\omega_2)Z_2Z_1^{*2}e^{i\theta t})/|Z_1|^2$. The reduction in the amplitude of $|Z_2|$ is equal to a half of $|Z_1|^2$.

## 4. DISCUSSION

Before closing, the necessary amplitude of the mother mode to actually observe the growth of daughter mode is discussed. Eq. (14) indicates that, if the damping rate becomes smaller (closer to the linear stability boundary), the threshold amplitude of mother mode for parametric instability becomes smaller. When the mother mode frequency is chirping, the frequency mismatch goes to zero as time goes. Thus, for the case of small $\gamma_d$, Eq. (14) can be satisfied for small value of $Z_{2,0}$ (when mother mode frequency is chirping as is in experimental observation). It looks that if $\gamma_d$ goes to zero, necessary value of $Z_{2,0}$ to satisfy Eq. (14) can approach to zero. However, for the very small values of $Z_{2,0}$ and $\theta$, the onset of daughter mode cannot be observed. Let us denote the value of mismatch $\theta$ by $\theta_c$, when marginal condition (from Eq. (14)) is satisfied. The time period $t_{duration}$, during which Eq. (14) is satisfied, is evaluated as

$$t_{duration} = 2\theta_c\left(\frac{d\theta}{dt}\right)^{-1}, \tag{25}$$

where $d\theta/dt$ is the frequency chirping rate of the mother mode. In order that the growth of the daughter mode is observed, the time period multiplied by the growth rate must be longer than $e$. We evaluate $\theta_c \sim 2\sqrt{D^2 - \gamma_d^2}$ and typical growth rate $\sim D - \gamma_d$, and the condition that the growth of the daughter mode is observed is written as

$$(D - \gamma_d)\sqrt{D^2 - \gamma_d^2} \gtrsim \frac{e}{4}\frac{d\theta}{dt} \sim \frac{d\theta}{dt}. \tag{26a}$$

In the small $\gamma_d$ limit, one has

$$D^2 \gtrsim \frac{d\theta}{dt}. \tag{26b}$$

Equations (26a) and (26b) describe the minimum amplitude for observing the growth of the daughter mode.

The 'abruptness of onset' is also reproduced in this model. We consider the case that the frequency mismatch $\theta$ is changing up in time owing to the frequency chirping of the mother mode, and that $d\theta/dt$ causes the temporal change of the growth rate (given by Eq. (13)). The time derivative of the growth rate, $\gamma = \mathrm{Re}(\lambda)$, is evaluated by taking into account of $d\theta/dt$. One has

$$\frac{\partial\gamma}{\partial t} = \frac{-(\theta/2)(\partial\theta/\partial t)}{2\sqrt{D^2 - \theta^2/4}}. \tag{27}$$

At the onset of daughter mode, the frequency mismatch satisfies the relation $\theta_c = 2\sqrt{D^2 - \gamma_d^2}$, so that

$$\frac{\partial\gamma}{\partial t} = \frac{\theta_c}{4\gamma_d}\left(-\frac{\partial\theta}{\partial t}\right) = \frac{\sqrt{D^2 - \gamma_d^2}}{2\gamma_d}\left|\frac{\partial\theta}{\partial t}\right|. \tag{28}$$

The characteristic growth rate of the daughter mode $\gamma_0$ is given as

$$\gamma_0 \sim D - \gamma_d. \tag{29}$$

The parameter that denotes the abruptness,

$$\frac{1}{\gamma_0^2}\frac{\partial\gamma}{\partial t} \tag{30}$$

(i.e., the onset of the mode is very fast if $\gamma_0^{-2}(\partial\gamma/\partial t) \gtrsim 1$), is given as

$$\frac{1}{\gamma_0^2}\frac{\partial\gamma}{\partial t} = \frac{\sqrt{D^2 - \gamma_d^2}}{2\gamma_d(D - \gamma_d)^2}\left|\frac{\partial\theta}{\partial t}\right| \sim \frac{1}{\gamma_d D}\left|\frac{\partial\theta}{\partial t}\right|. \tag{31}$$

We see that the change of the growth rate is governed by the chirping rate of the mother mode (which is owing to the nonlinear dynamics in the wave response), not by the change of global plasma parameters. In the experimental observation [7], parameters were evaluated as $-\partial\theta/\partial t \sim 5\times10^{-4}\omega_1^2$, $\gamma_d \sim 0.03\omega_1$, and $D \gtrsim \gamma_d$ [13]. Substituting these parameters into Eq. (31), the estimate is given as

$$\frac{1}{\gamma_0^2}\frac{\partial\gamma}{\partial t} \sim O(1). \tag{32}$$

We see that the present model can explain the abruptness of the onset of daughter mode, which is observed in experiments [7].

At the end, a possible further application of this model is commented. This model can be applied to the case of Alfvén eigenmodes, which have been known to evolve nonlinearly with frequency chirping. In addition, they could become subcritically unstable. Thus, the model in this article could be applied to the case of Alfvén eigenmodes. The application to neoclassical tearing mode, which is known to be subcritically unstable, also can be possible. In the analysis [21], it was demonstrated that the ambient turbulence transfers the energy so as to excite the neoclassical tearing mode. The parametric coupling of neoclassical tearing mode with chirping mode merits further studies.

## 5. SUMMARY

In this article, we discussed the analytic model for the abrupt onset of GAM in the presence of chirping EGAM. Threshold conditions for the onset of abrupt growth are derived, and the period doubling phenomenon is explained. The phase relation between the mother and daughter modes are also discussed. It explains that the onset of daughter mode takes place when the frequency of the mother mode is slightly lower than the twice of the daughter mode frequency. The phenomenon is also interpreted as a period doubling bifurcation of the mother mode (EGAM). It is noticed that conditions of Eqs. (14) and (26a) are close to what have been empirically observed by direct nonlinear simulations (Fig. 3 of [13]). The increasing rate of the growth rate is discussed. It was shown that the growth rate reaches the characteristic large value $\gamma_0$ within the short time, which is of the order of $1/\gamma_0$. Thus, the onset condition as well as the abruptness of the excitation is derived in an analytic manner. They are the essential elements in the 'trigger problem', which has been studied in conjunction with the onset of large-scale deformation in the wide area of plasma physics. The result helps as a guide for experiments. The analysis here contributes to the understanding of 'trigger problems' of plasmas in laboratory and nature.

Discussion with S. Inagaki and K. Ida is acknowledged. This work was supported by Grant-in-Aid for Scientific Research (15H02155, 23244113), NIFS/NINS under NIFS10ULHH020, and the Collaborative Research Programs of Research Institute for Applied Mechanics, Kyushu University and of NIFS, and by Asada Science Foundation.

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