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<td>Author(s)</td>
<td>Omura, Yoshiharu; Nunn, David</td>
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<tr>
<td>Citation</td>
<td>Journal of Geophysical Research (2011), 116</td>
</tr>
<tr>
<td>Issue Date</td>
<td>2011-06</td>
</tr>
<tr>
<td>URL</td>
<td><a href="http://hdl.handle.net/2433/141944">http://hdl.handle.net/2433/141944</a></td>
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<td>Rights</td>
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<td>Type</td>
<td>Journal Article</td>
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Kyoto University
Triggering process of whistler-mode chorus emissions in the magnetosphere

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Abstract. Chorus emissions are triggered from the linear cyclotron instability driven by the temperature anisotropy of energetic electrons (10 - 100 keV) in the magnetosphere. Chorus emissions grow as an absolute nonlinear instability near the magnetic equator due to the presence of an electromagnetic electron hole in velocity space. The transition process from the linear wave growth at a constant frequency to the nonlinear wave growth with a rising tone frequency is due to formation of a resonant current $- J_B$ anti-parallel to the wave magnetic field. The rising-tone frequency introduces a phase shift to the electron hole at the equator, and results in a resonant current component anti-parallel to the wave electric field $- J_E$, which causes the nonlinear wave growth. To confirm this triggering mechanism, we perform Vlasov Hybrid Simulations with $J_B$ and without $J_B$. The run without $J_B$ does not reproduce chorus emissions, while the run with $J_B$ does successfully reproduce chorus emissions. The nonlinear frequency shift $\omega_1$ due to $J_B$ plays a critical role in the triggering process. The nonlinear transition time $T_N$ for the frequency shift is found to be of the same order as the nonlinear trapping period, which is confirmed by simulations and observation. The established frequency sweep rate is $\omega_1/T_N$, which gives an optimum wave amplitude of chorus emissions.
1. Introduction

Coherent electromagnetic waves called chorus emissions have frequently been observed in the inner magnetosphere [e.g., Tsurutani and Smith, 1974; Santolik et al., 2003; Santolik, 2008; Kasahara et al., 2009]. Chorus emissions typically consist of a series of rising tones generated near the magnetic equator, excited by energetic electrons from several keV to tens of keV injected into the inner magnetosphere at the time of a geomagnetic disturbance. In recent years chorus emissions have been studied extensively because of their role as a viable mechanism for accelerating radiation belt electrons as well as precipitating them into the polar atmosphere [Summers et al., 1998; Summers et al., 2002; Miyoshi et al., 2003; Horne et al., 2005; Omura and Summers, 2006; Omura et al., 2007; Katoh and Omura, 2007; Furuya et al., 2008; Katoh et al., 2008; Hikishima et al., 2010].

As a generation mechanism for chorus emissions, a backward wave oscillator (BWO) theory has been proposed by Trakhtengerts [1999]. The BWO theory assumes a step in the velocity distribution function, which has not been observed. Furthermore such a step can only influence the 'triggering phase' and can have no effect upon the plasma dynamics of the established chorus element with sweeping frequency as the resonance velocity will no longer match the location of the step in parallel velocity.

Numerical modeling of chorus emissions have been performed using a Vlasov-Hybrid Simulation (VHS) code based on narrow band field equations derived from Maxwell’s equations and the linear equation of motion of cold plasma, under the assumption of a band-limited coherent whistler-mode wave [Nunn, 1990, 1993; Nunn et al., 1997, 2009]. The initial wave amplitude and the wave phase are specified in such simulations. In
contrast to the VHS code, chorus emissions with rising tones were reproduced successfully
starting from thermal noise in an electromagnetic electron-hybrid code, in which Maxwell’s
equations are solved directly together with the electron fluid equation for the cold dense
electrons and the equations of motion for the hot resonant electrons [Katoh and Omura,
2006; 2007]. The mechanism of the rising chorus emissions has been analyzed theoretically
in terms of nonlinear wave growth due to the formation of an electromagnetic electron
hole in the velocity phase space [Omura et al., 2008, 2009]. The relation between the wave
amplitude and the frequency sweep rate in the generation region of chorus emissions has
been derived [Omura et al., 2008, Equation (50)]. The validity of this relation has been
demonstrated in a full-particle electromagnetic simulation [Hikishima et al., 2009] as well
as in the electron-hybrid simulation [Katoh and Omura, in press]. These simulations show
that seeds of chorus emissions with rising tones are formed in a localized region near the
magnetic equator. The seeds of emissions grow as a result of the formation of resonant
current arising from nonlinear trajectories of resonant untrapped electrons.

Falling tone chorus emissions have also been observed by Cluster spacecraft, and mod-
eled by Nunn et al. [2009]. In the present analysis, however, we focus our attention on
rising tone chorus emissions, which has been studied more extensively by simulations [Ka-
toh and Omura, 2007; Hikishima et al., 2009; Nunn et al., 2009] as well as observations
[Macusova et al., 2010; Cully et al., 2011].

The nonlinear wave growth theory [Omura et al., 2008, 2009] assumed a frequency sweep
rate \( \partial \omega / \partial t \) that drives the wave growth through the formation of a resonant current \( J_E \)
parallel to the wave electric field. However, the mechanism for the formation of finite
\( \partial \omega / \partial t \) has not been clarified yet. Noting that that the resonant current \( J_B \) parallel to the
wave magnetic field was included in previous studies based on the VHS code [Nunn et al., 1997, 2009], we have conducted an experiment to run the VHS code with and without $J_B$ to find out the contribution of $J_B$ to the generation of chorus emissions. As shown in Figure 1, the run without $J_B$ does not reproduce a rising tone chorus element, but a constant frequency emission, while the run with $J_B$ does. Considering the contribution of resonant electrons, we first analyze the optimum condition for triggering chorus theoretically in section 2. We confirm the theoretical model by analyzing the result of the Vlasov Hybrid Simulation in section 3. In section 4 we present a summary and discussion.

2. Condition for triggering chorus emissions

We assume a whistler-mode wave propagating parallel to the static magnetic field $B_0$ with a wavenumber $k$ and a constant frequency $\omega_0$ satisfying the linear dispersion relation

$$c^2 k^2 - \omega_0^2 - \frac{\omega_0 \omega_{pe}^2}{\Omega_e - \omega_0} = 0.$$ (1)

where $c$, $\Omega_e$, and $\omega_{pe}$ are the speed of light, the electron cyclotron frequency, and the electron plasma frequency, respectively. It is noted that the wavenumber $k$ is a function of a distance $h$ taken along the magnetic field line from the magnetic equator. The wave fields are in the transverse plane containing $x$- and $y$-axes. We then assume energetic electrons interacting with the wave satisfying the cyclotron resonance condition

$$\omega_0 - kv_\parallel = \frac{\Omega_e}{\gamma},$$ (2)

where $\gamma$ is the Lorentz factor given by $\gamma = [1 - (v_\parallel^2 + v_\perp^2)/c^2]^{-1/2}$, and $v_\parallel$ and $v_\perp$ are electron velocities parallel and perpendicular to $B_0$. Solving for $v_\parallel$, we can obtain an
explicit expression of the cyclotron resonance velocity

\[ \tilde{V}_R = \frac{\tilde{\omega}^2 - \sqrt{\tilde{\omega}^4 + (\tilde{\omega}^2 + \tilde{V}_p^2)(1 - \tilde{\omega}^2 - \tilde{V}_{\parallel 0}^2)}}{\tilde{\omega}^2 + \tilde{V}_p^2} \tilde{V}_p, \quad (3) \]

where \( \tilde{\omega} = \omega_0/\Omega_e, \tilde{V}_R = V_R/c, \tilde{V}_p = V_p/c, \) and \( \tilde{V}_{\parallel 0} = V_{\parallel 0}/c. \) The phase velocity is given by

\[ V_p = \frac{\omega_0}{k} = c\delta \xi, \quad (4) \]

where \( \xi^2 = \omega_0(\Omega_e - \omega_0)/\omega_{pe}^2 \) and \( \delta^2 = (1 + \xi^2)^{-1}[Omura et al., 2008]. \) When we evaluate \( \gamma \) in the equations derived below, we substitute \( v_\parallel = V_R \) and \( v_\perp = V_{\perp 0}, \) where \( V_{\perp 0} \) is the average perpendicular velocity.

The electrons are organized in phase in the transverse plane, and form a resonant current that can trigger formation of a new wave field with a variable frequency \( \omega, \) as demonstrated by a full-particle simulation by Hikishima et al. [2010].

We express the electric and magnetic field vectors of the total wave field in the transverse plane by the complex forms \( \tilde{E}_w = E_w \exp(i\psi_E) \) and \( \tilde{B}_w = B_w \exp(i\psi_B), \) respectively.

From Maxwell’s equations we obtain the following equations for the amplitude \( B_w \) of the wave magnetic field in the form [Omura et al., 2008],

\[ \frac{\partial B_w}{\partial t} + V_g \frac{\partial B_w}{\partial h} = -\frac{\mu_0 V_g}{2} J_E, \quad (5) \]

\[ c^2 k^2 - \omega^2 - \frac{\omega \omega_{pe}^2}{\Omega_e - \omega} = \mu_0 c^2 k \frac{J_B}{B_w}, \quad (6) \]

where \( \mu_0 \) is the magnetic permittivity in vacuum. The resonant current formed by resonant electrons is divided into two components \( J_E \) and \( J_B \) parallel to the transverse wave electric and magnetic fields, respectively. Details of the derivation of (5) and (6) are found in Appendix A of Omura et al.[2008].
While the resonant current $J_E$ modifies the wave amplitude $B_w$, the quantity $J_B/B_w$ changes the frequency $\omega$ of the triggered wave. It is noted that the wavenumber $k$ or the wavelength does not change in space and time because it is imposed by the triggering wave with the constant frequency $\omega_0$ in the present situation. Denoting the frequency deviation from $\omega_0$ as $\omega_1 (\omega = \omega_0 + \omega_1)$ and assuming $\omega_1 << \omega_0$, we expand (6) around $\omega_0$ to obtain

$$\left\{ 2\omega_0 + \frac{\Omega_e \omega_{pe}^2}{(\Omega_e - \omega_0)^2} \right\} \omega_1 = -\mu_0 c^2 k \frac{J_B}{B_w} , \quad (7)$$

where we made use of (1). Differentiating (1) with respect to $\omega_0$, we obtain

$$2e^2 k \frac{\partial k}{\partial \omega_0} = 2\omega_0 + \frac{\Omega_e \omega_{pe}^2}{(\Omega_e - \omega_0)^2} . \quad (8)$$

Using (8), we rewrite (7) as

$$\omega_1 = -\frac{\mu_0 V_g}{2} \frac{J_B}{B_w} , \quad (9)$$

where $V_g = \partial \omega_0 / \partial k$.

We consider the optimum condition for the nonlinear wave growth to take place, assuming that the electron hole is progressively formed in the velocity phase space within the time window defined by the nonlinear transit time $T_N$. Dynamics of a resonant electron is described by a set of simplified differential equations. Introducing the variables

$\theta = k(v_\parallel - V_R)$ and the phase angle $\zeta$ between the perpendicular velocity $v_\perp$ and the wave magnetic field $B_w$,

$$\frac{d\zeta}{dt} = \theta , \quad (10)$$

and

$$\frac{d\theta}{dt} = \omega_{tr}^2 (\sin \zeta + S) , \quad (11)$$
where \( \omega_{tr} = \omega_t \delta \gamma^{-1/2} \), and we have assumed that \( v_\parallel \sim V_R \), i.e., \( \theta \sim 0 \). Here, \( \omega_t \) is the trapping frequency given by \( \omega_t = \sqrt{k v_\perp \Omega_w} \). The shape of the electron hole is determined by the inhomogeneity ratio given by

\[
S = -\frac{1}{s_0 \omega_w} \left( s_1 \frac{\partial \omega}{\partial t} + c s_2 \frac{\partial \Omega_e}{\partial h} \right), \tag{12}
\]

where

\[
s_0 = \frac{\delta V_{\perp 0}}{\xi c}, \tag{13}
\]

\[
s_1 = \gamma (1 - \frac{V_R}{V_g})^2, \tag{14}
\]

and

\[
s_2 = \frac{1}{2 \xi \delta} \left\{ \frac{\gamma \omega_0}{\Omega_e} \left( \frac{V_{\perp 0}}{c} \right)^2 - \left[ 2 + \Lambda \frac{\delta^2 (\Omega_e - \gamma \omega_0)}{\Omega_e - \omega_0} \right] \frac{V_R V_p}{c^2} \right\}. \tag{15}
\]

We have incorporated the variation of the cold electron density \( N_e(h) \) along the magnetic field line as \( N_e(h) = N_{e0} \Omega_e(h)/\Omega_{e0} \), where \( N_{e0} \) and \( \Omega_{e0} \) are respectively the cold electron density and the electron gyrofrequency at the equator. We have \( \Lambda = \omega_0/\Omega_e \) for this inhomogeneous electron density model, while \( \Lambda = 1 \) for the constant electron density model [Omura et al. 2009].

The resonant currents \( J_E \) and \( J_B \) are expressed respectively as

\[
J_E = -J_0 \int_{\zeta_1}^{\zeta_2} \left[ \cos \zeta_1 - \cos \zeta + S(\zeta - \zeta_1) \right]^{1/2} \sin \zeta d\zeta, \tag{16}
\]

and

\[
J_B = J_0 \int_{\zeta_1}^{\zeta_2} \left[ \cos \zeta_1 - \cos \zeta + S(\zeta - \zeta_1) \right]^{1/2} \cos \zeta d\zeta, \tag{17}
\]

where \( J_0 = (2e)^{3/2}(m_0 k \gamma)^{-1/2} V_{\perp 0}^{5/2} \delta Q GB_w^{1/2} \), and \( e \) and \( m_0 \) are the charge and rest mass of an electron. The factor \( Q \) represent the depth of the electron hole [Omura et al., 2009].
The phase angles $\zeta_1$ and $\zeta_2$ define the boundary of the trapping wave potential as described by Omura et al. [2009]. The parameter $G$ is the value of the velocity distribution function $g(v_\parallel,\zeta)$ in the trapping region around the resonance velocity.

We assume that the velocity distribution function $f$ of hot energetic electrons is given in terms of the relativistic momentum per unit mass $u = \gamma v$; $u$ has components $u_\parallel = \gamma v_\parallel$ and $u_\perp = \gamma v_\perp$, respectively parallel and perpendicular to the ambient magnetic field. We specify $f$ as

$$f(u_\parallel, u_\perp) = \frac{N_h}{(2\pi)^{3/2} U_\perp U_\parallel 0} \exp \left( -\frac{u_\parallel^2}{2U_\parallel^2} \right) \Delta(u_\perp - U_\perp 0) ,$$

where $U_\perp 0 = \gamma V_\perp 0$, $U_\parallel$ is the thermal momentum in the parallel direction, and $\Delta$ is the Dirac delta function, and we have normalized $f$ to the density of hot electrons $N_h$.

Integrating over $u_\perp$ and taking an average over $\zeta$, we obtain the magnitude $G$ of the unperturbed distribution function $g(v_\parallel,\zeta)$ at the resonance velocity $V_R$ as

$$G = \frac{N_h}{(2\pi)^{3/2} U_\parallel 0} \exp \left( -\frac{\gamma^2 V_R^2}{2U_\parallel^2} \right) ,$$

Omura et al. [2008] found that the maximum value of $-J_E$ takes place when $S = -0.413$.

Solving (17) for $S = 0.413$, we obtain $J_B = -1.3J_0$, which is rewritten as

$$J_B = -1.3(2e)^{3/2} \left( \frac{B_w}{m_0 k \gamma} \right)^{1/2} V_\perp 0^{5/2} \delta Q G .$$

The nonlinear transition time $T_N$ for formation of the nonlinear resonant current is roughly estimated by the nonlinear trapping period $T_tr$ given by

$$T_tr = \frac{2\pi}{\omega_tr} = \frac{2\pi}{\delta} \left( \frac{m_0 \gamma}{k V_\perp 0 e B_w} \right)^{1/2} ,$$

where $\omega_tr$ is the trapping frequency [Omura et al., 2008]. We define a ratio $\tau = T_N/T_tr$, which is to be determined by numerical simulations in the next section.
Through the nonlinear transition time $T_N$, the electron hole is gradually formed. Along with formation of $J_B$, the frequency of the triggered wave gradually deviates from $\omega_0$ to $\omega_0 + \omega_1$. From (9), (19), (20), and (21), we obtain the frequency sweep rate over the nonlinear trapping period as

$$\frac{\omega_1}{T_N} = \frac{1.3}{4} \pi^{-5/2} Q \left( \frac{\omega_{ph} V_{\perp 0} \delta}{\gamma c} \right)^2 \frac{V_g}{U_{t\parallel}} \exp \left( -\frac{\gamma^2 V_R^2}{2 U_{t\parallel}^2} \right),$$

(22)

where $\omega_{ph}$ is the plasma frequency of the hot energetic electrons defined by $\omega_{ph}^2 = \mu_0 c^2 N_h e^2/m_0$.

At the equator the inhomogeneity of the magnetic field is zero, and the second term on the right-hand side of (12) vanishes. Since the maximum nonlinear wave growth takes place when $S = -0.4$ [Omura et al., 2008], we can derive from (12) the relation between the frequency sweep rate and the normalized wave amplitude at the equator $\Omega_w = eB_w/m_0$ in the form,

$$\frac{\partial \omega}{\partial t} = \frac{0.4 s_0 \omega_0}{s_1} \Omega_w .$$

(23)

Equating the left-hand sides of (22) and (23), we obtain an optimum wave amplitude $\Omega_{wo}$ that can trigger the rising-tone chorus element as

$$\Omega_{wo} = \frac{s_1 \omega_1}{0.4 s_0 \omega_0 T_N} .$$

(24)

From (22) and (24), we obtain

$$\tilde{\Omega}_{wo} = 0.81 \pi^{-5/2} \frac{Q}{\tau \ s_0 \omega \bar{U}_{t\parallel} \left( \frac{\bar{\omega}_{ph} \bar{V}_{\perp 0} \delta}{\gamma} \right)} \exp \left( -\frac{\gamma^2 \bar{V}_R^2}{2 \bar{U}_{t\parallel}^2} \right),$$

(25)

where $\tilde{\Omega}_{wo} = \Omega_{wo}/\Omega_{e0}$, $\bar{\omega}_{ph} = \omega_{ph}/\Omega_{e0}$, and $\bar{U}_{t\parallel} = U_{t\parallel}/c$. 
Finally we can evaluate the nonlinear transition time using the wave amplitude obtained above. Using (4), we rewrite (21) as

\[ T_N \Omega_{e0} = 2\pi \tau \left( \frac{\gamma \xi}{\tilde{\omega} \tilde{V}_{\perp0} \Omega_{e0} \delta} \right)^{1/2} , \] (26)

The triggered chorus element should satisfy another condition to grow as a nonlinear absolute instability at the equator. The wave amplitude of the triggered emission should be greater than the threshold for the nonlinear wave growth \cite{Omura et al., 2009}, which is given by

\[ \tilde{\Omega}_{th} = 1000 \pi^3 \gamma^3 \xi \left( \frac{\tilde{a}s_2 \tilde{U}_{t\|}}{Q} \right)^2 \exp \left( \frac{\gamma^2 \tilde{V}_R^2}{\tilde{U}_{t\|}^2} \right) , \] (27)

where \( \tilde{a} = \frac{ac^2}{\Omega_{e0}^2} \), and \( a \) is a coefficient defining the parabolic variation of the magnetic field around the equator as \( \Omega_e = \Omega_{e0}(1 + ah^2) \). The coefficient \( a \) is specified by the \( L \) value and the Earth’s radius \( R_E \) as \( a = 4.5/(LR_E)^2 \).

3. Simulations of rising tone emissions

We performed simulations of rising tone emissions by the VHS code, in which the equations equivalent to (5) and (6) are solved numerically along with the Vlasov equation for the resonant electrons \cite{Nunn et al., 1990, 1993, 2009}. To find out the importance of the resonant current \( J_B \) in generating chorus emissions, we first performed a run in which both \( J_E \) and \( J_B \) are calculated and the wave field is updated by them. The run produced a rising tone emission as shown in Figure 1(a). We then performed the second run with \( J_B = 0 \). The second run did not reproduce the rising tone emission as shown in Figure 1(b). A constant frequency emission did result with symmetric upper and lower resonant sidebands. We now recognize the importance of the resonant current \( J_B \) as analyzed in the preceding section.
The parameters used in the VHS run are the following. The electron cyclotron frequency $f_c$ is 8 kHz at the magnetic equator. The electron plasma frequency $f_p$ is 18.8 kHz, which gives $\tilde{\omega}_p = 2.35$. Assuming $L = 4.79$, we have $\tilde{a} = 1.72 \times 10^{-7}$. The velocity distribution assumed in the VHS code is approximated by (18) with $V_{\perp 0} = 0.4584c$, $U_{t\parallel} = 0.2396c$, and $\tilde{\omega}_{ph} = 0.1664 \left( N_h/N_e = 0.005 \right)$. These parameter gives equatorial linear growth rate of 600 dB/s, which is in rough agreement with path integrated gain of whistler-mode waves evaluated based on the THEMIS spacecraft observation [Li et al., 2009].

Figures 2(a) and 2(b) show the spatial and temporal variation, in dimensionless units, of wave amplitude $B_w/B_0$ and wave frequency $\omega/\Omega_{e0}$, respectively. The incoming triggering pulse has an initial amplitude $B_{w0} = 4.2 \times 10^{-5}$ (12pT) and duration $8.4 \times 10^3 \Omega_{e0}^{-1}$ (167 ms) and frequency $\omega/\Omega_{e0} = 0.4$ (3.2 kHz). Its amplification as it progresses to the equator is clearly visible in Figure 2a, at which point its amplitude is large enough for nonlinear trapping to occur and for triggering to take place. At about $t\Omega_{e0} \sim 9000$ a riser generation region is firmly established with a wave generation point at the equator and a profile leading edge some $100 \sim 500$ units upstream from the equator.

From the temporal and spatial variation of local frequency in Figure 2(b) we see an initial drop in frequency due to nonlinear trapping in the incoming pulse when it is in the positive inhomogeneity region upstream from the equator $S > 0$. The region of much lower frequencies near $t\Omega_{e0} \sim 8000$ is not significant as it corresponds to very small amplitudes in the wake of the incoming pulse. After $t\Omega_{e0} \sim 9000$ we see a rising frequency in the equatorial region, and after $t\Omega_{e0} \sim 11000$ the progressive establishment of a frequency gradient across the interaction region.
By plotting the distribution function of resonant electrons in the \( (v_{\parallel} - \zeta) \) phase space, we find that the depletion of the electrons in the trap at the resonance velocity is about 22 percent of the surrounding energetic electrons. Assuming \( Q = 0.25 \) with the parameters of the VHS run, we calculated the optimum wave amplitudes given by (25) in solid lines in Figure 3(a). The numbers attached to the lines are the time scale factors \( \tau \). We also plot the threshold for the nonlinear wave growth given by (27) in a dashed line. Figure 3(b) shows the nonlinear time scale \( T_N \) given by (26) for different values of \( \tau \).

In Figure 4 we plot simultaneous time histories at the equator \( (h = 0) \) of wave amplitude and frequency as well as resonant particle current components \( J_E \) and \( J_B \) in arbitrary units. Figure 4(a) shows wave amplitude progression. Significant amplitudes \( \sim 2 \times 10^{-4} \) are established after \( t\Omega_{e0} = 5000 \), but there is a drop out around \( t\Omega_{e0} \sim 9000 \) when the wave profile slips downstream from the equator. Figure 4(b) shows the progression of equatorial frequency. After an initial drop a positive frequency gradient is abruptly established at \( t\Omega_{e0} \sim 9000 \) of magnitude \( 1 \times 10^{-5}\Omega_{e0}^2 \) (4 kHz/s). From (23), we can estimate the average wave amplitude \( B_w/B_{oEQ} = 1.3 \times 10^{-4} \) (37 pT) that results in the frequency sweep rate. Assuming this amplitude is the optimum amplitude given by (25), we find \( \tau = 1.0 \). Substituting this value of \( \tau \) into (26), we find \( T_N\Omega_{e0} = 640 \) (13 ms).

Figure 4(e) shows the time development of the nonlinear frequency shift term \( \omega_1 \). The latter is quite interesting. At around \( t\Omega_{e0} \sim 9500 \) it rises quite quickly to a value \( \sim 0.006 \) (50 Hz) which is then sustained. The development time is nearly equal to \( T_N \) obtained above. We can also confirm that the quotient \( \omega_1/T_N \) agrees with the frequency sweep rate \( \partial\omega/\partial t = 1 \times 10^{-5}\Omega_{e0}^2 \) in accordance with the assumption of the optimum wave amplitude.
As a result of trapping in the negative inhomogeneity region $h > 0$, we get an electron hole and a significant build up of a large nonlinear current $J_B$ at the equator as shown in Figure 4(c). This is as expected since the inhomogeneity $S$ at $h = 0$ is initially zero giving a phase trapping angle anti parallel to the wave magnetic field (phase of 180 degrees relative to $B_w$). We find formation of a negative $J_B$ around $t\Omega_{e0} = 6500 \sim 8500$, but the wave amplitude is much larger than the optimum wave amplitude, and the nonlinear frequency shift was too small to trigger the rising tone. As the positive frequency sweep rate becomes established at later time as shown in Figure 4(b), $S$ shifts to a value $\sim -0.4$ which enables us to easily interpret the $J_E$ profile in Figure 4(d). The current $J_E$ is initially close to zero as expected for $S = 0$, but acquires a significant negative component, from about $t\Omega_{e0} \sim 9500$ giving nonlinear growth, as the trapping angle rotates to a phase of approximately 166 degrees.

We performed a few other runs of rising chorus elements by the VHS code with different parameters, and found the same sequence of the initial formation of the negative $J_B$ followed by gradual formation of the negative $J_E$ along with establishment of the finite frequency sweep rate. We also found $\tau \sim 1.0$.

In Figure 5(a) we plot the optimum wave amplitude and the threshold for the nonlinear wave growth with parameters used in an electron hybrid simulation [Katoh and Omura [2007]: Omura et al., 2008] $\tilde{a} = 9.8 \times 10^{-7}, \tilde{V}_{\perp0} = 0.6, \tilde{U}_{||} = 0.3, \tilde{\omega}_{pe} = 4$, and $\tilde{\omega}_{ph} = 0.113$. We assume $Q = 0.5$ and different values of $\tau = 0.25, 0.5, 1.0, 2.0$. Noting that the simulated chorus emissions have frequency spectra starting from $\omega = 0.2\Omega_{e0}$ with the wave amplitude $B_w \sim 4 \times 10^{-4}B_{0EQ}$ at the equator (Figures 4 and 5 of Omura et al. [2008]), we find that the optimum amplitude with $\tau = 0.25 \sim 0.5$ agrees with the simulation result, because
the frequency rage above $\omega = 0.2$ satisfies the necessary condition for the nonlinear wave
growth $\tilde{\Omega}_{wo} > \tilde{\Omega}_{th}$.

In Figure 5(b), we plot the optimum wave amplitude and the threshold with parameters
used in a full-particle simulation by Hikishima et al. [2009]: $\tilde{a} = 5.1 \times 10^{-6}$, $\tilde{V}_{\perp 0} = 0.29$,
$\tilde{U}_{\| l} = 0.2$, $\tilde{\omega}_{pe} = 5$, and $\tilde{\omega}_{ph} = 0.40$. In both cases, we assume $Q = 0.5$ and different
values of $\tau = 0.25, 0.5, 1.0, 2.0$. Figure 8 of Hikishima et al. [2009] shows chorus emissions
starting from $\omega = 0.16\Omega_{eo}$ with a wave amplitude $\tilde{B}_w \sim 1 \times 10^{-3}\tilde{B}_0EQ$ at the equator.

We find good agreement between the optimum amplitude with $\tau = 0.5$ and the frequency
spectra of the simulation result.

It is also interesting to note that the frequency range that satisfies $\tilde{\Omega}_{wo} > \tilde{\Omega}_{th}$, shown
in Figure 5(a), agrees with the width of the wave spectra found in the simulation result.

A chorus element undergoes a strong wave amplitude modulation through its evolution
in frequency. Even if the nonlinear wave growth is terminated because of the amplitude
modulation, a new triggering process can take place when the optimum condition of the
wave amplitude is satisfied, and the rising tone is resumed. Therefore, while the optimum
wave amplitude is larger than the threshold for the nonlinear wave growth, the frequency
continues to increase, forming the chorus element.

4. Summary and Discussion

We have obtained an optimum wave amplitude that can trigger a rising tone chorus
element. When the optimum wave amplitude given by (25) is reached by the linear
wave instability at a specific frequency driven by the temperature anisotropy of resonant
electrons near the equator, a triggered emission arises with a rising frequency due to
gradual formation of the negative $J_B$ and with an increasing wave amplitude due to the
negative $J_E$ induced by the positive frequency sweep rate. An electromagnetic electron
hole is formed in the velocity phase space with $S \sim -0.4$. An absolute nonlinear instability
takes place above the threshold given by (27).

For comparison with an observation by the Cluster spacecraft [Santolik et al., 2003;
Santolik, 2008], we plot the optimum wave amplitude and the threshold for the nonlinear
wave growth in Figure 6(a), and the theoretical frequency sweep rate $\partial f / \partial t$ in Figure 6(b).

Based on the observation at $L = 4.4$, which gives $a = 2 \times 10^{-7}$, we assumed parameters
as $\omega_{pe}/\Omega_{e0} = 2.4$, $V_{\perp 0} = 0.3$, and $U_{t\parallel} = 0.2$. The energy of the resonant electrons are
50 keV for $\omega/\Omega_{e0} = 0.37$ (3 kHz). Since the parameters are very close to the VHS run,
we use $Q = 0.25$ and varied the time scale parameter as $\tau = 0.25, 0.5, 1.0, 2.0$. The
optimum wave amplitude and the frequency sweep rate varies depending on the density
of energetic electrons $N_h$. From the observation, we can find the frequency sweep rates,
which do not change much through propagation. On the other hand, the wave amplitude
can change substantially through propagation from the source region to the observation
point. Therefore, we can infer the physical parameters at the generation region from
the frequency sweep rate of the observed chorus emissions. Chorus emissions reported
by Santolik et al.[2003] consist of rising tone elements starting from $2 \sim 3$ kHz with the
frequency sweep rate $\sim 15$ kHz/s. We find a good agreement with the case $N_h/N_c = 0.04$
and $\tau = 0.25 \sim 0.5$. The wave amplitude in the generation region is estimated as large as
$100 \sim 300$ pT.

It is also interesting to note that the frequency sweep rate give by (22) does not depend
on the wave amplitude of the triggering wave. The physical plasma parameters and the
frequency of the triggering wave determine the frequency sweep rate. Then the effective
wave growth takes place with the wave amplitude of the triggered wave, which satisfies the condition (23) at the magnetic equator.

Finally we can construct the following scenario of the chorus generation process, summarizing the recent studies [Omura et al., 2008, 2009] and the present analysis.

1. **Linear Phase**: Whistler-mode waves grow due to a linear instability driven by temperature anisotropy of energetic electrons. A constant-frequency wave with the maximum linear growth rate becomes dominant to form a coherent wave phase. The linear growth rate maximizes at the magnetic equator where the flux of energetic electrons at the resonance velocity becomes largest. The seed of the triggering wave is the thermal fluctuation at some distance from the equator. It propagates toward the equatorial region, undergoing the convective linear growth to form a coherent wave near the equator.

2. **Nonlinear Phase 1**: As a result of the nonlinear dynamics of electrons at or near the resonance velocity of the triggering coherent wave, a hole in the velocity distribution function at the location of the resonant particle trap is formed. The negative resonant current $J_B$ is formed, resulting in the nonlinear frequency shift $\omega_1$ taking place over $T_N$, which is of the same order as the nonlinear trapping time $T_{tr}$. The ratio $T_N/T_{tr}$ ($= \tau$) is $0.25 \sim 1.0$. The frequency sweep rate is the nonlinear frequency shift $\omega_1$ divided by the nonlinear transition time $T_N$.

3. **Nonlinear Phase 2**: Along with the establishment of the frequency sweep rate $\omega_1/T_N$, the nonlinear resonant current $J_E$ is formed, resulting in the nonlinear wave growth. Formation of the negative $J_E$ results in the nonlinear wave growth as an absolute instability at the magnetic equator. The nonlinear instability starts from an optimum amplitude that is required to maximize $J_E$ with the inhomogeneity ratio $S = -0.4$. The
initial evolution of the wave amplitude and frequency is described by the chorus equations 
[Omura et al., 2009].

4. **Nonlinear Phase 3**: A chorus element is formed at the magnetic equator, and 
propagates away from it. As the wave propagates away from the equator, the zero order 
field gradient increases as does the wave amplitude due to the nonlinear wave growth. 
While $-1 < S < 0$, trapping is allowed and the chorus element continues to grow as it 
propagates away from the equator.

In the linear phase, the triggering wave grows from the thermal fluctuation through the 
convective linear growth. Evaluation of the path-integrated gain of the linear instability 
was performed by Li et al. [2009]. It is noted, however, that the linear growth rate 
is applicable only to the triggering waves, and not to the chorus emissions, which are 
generated through the nonlinear processes as described above.

The initial appearance of $J_B$ followed by $J_E$ in the initial nonlinear growth phase of the 
triggered emissions is also observed in the HAARP HF ionospheric heating experiment 
[Golkowski et al., 2010]. From the growth of the observed wave amplitudes and phases 
of HAARP-generated whistler-mode echoes, the magnitudes of the resonant currents are 
estimated. The analysis shows that the magnitude of $J_B$ is greater than that of $J_E$, which 
is consistent with the VHS code result shown in Figure 4 and the theoretical model of an 
electromagnetic electron hole assumed in the nonlinear wave growth theory.

Gibby et al. [2008] tried to reproduce triggered emissions, using a more classical Vlasov 
method such as developed by Denavit [1972; 1985] and by Besse and Sonnendrucker 
[2003] in the semi Lagrangian method. These methods advance the distribution function 
by applying Liouville’s theorem to one step particle trajectories, and then interpolating
distribution function back onto the phase space grid. These methods can have poor robustness to distribution function filamentation and give rise to an unphysical diffusion in distribution function due to the successive interpolation procedures. The VHS method suffers from neither of these problems. The code by Gibby et al. [2008] does not update wave phase properly when wave amplitude is very small and for a broadband simulation has a rather low resolution in phase space, which may be why the code does not trigger emissions. Gibby et al. [2008] presents interesting and plausible data, supported by his simulations, suggesting that in the key down case saturation may arise from marked spectral broadening which destroys particle trapping and thus nonlinear growth rates.

At the magnetic equator, the nonlinear wave growth saturates due to a subsequent nonlinear effect [Hikishima et al., 2010]. The detailed analysis of the saturation mechanism is left as a future study.

Acknowledgments. Computation in the present study was performed with the KDK system of Research Institute for Sustainable Humanosphere (RISH) at Kyoto University. This work was supported by Grant-in-Aid 20340135 of the Ministry of Education, Science, Sports and Culture of Japan. D. Nunn thanks Kyoto University for Visiting Professorship at RISH.

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**Figure 1.** Dynamic spectra of VHS code runs (a) with $J_B$ and (b) without $J_B$.

**Figure 2.** Spatial and temporal evolutions of (a) wave amplitude and (b) frequency in the VHS code run with $J_B$. The dashed white lines indicate the time when the formation of the nonlinear current $J_B$ begins. The arrow in magenta indicates the wave packet of the triggering wave.

**Figure 3.** (a) The optimum wave amplitudes (solid lines) for triggering rising tone emissions with different values of the time scale factor $\tau$ (attached numbers), and the threshold of the wave amplitude for the nonlinear wave growth (dashed line) with the parameters used in the VHS run. (b) The corresponding nonlinear transition time $T_N$ for formation of the nonlinear resonant current $-J_B$ with different values of the time scale factor $\tau$.

**Figure 4.** Time histories of the wave amplitude, frequency, resonant currents $J_B$ and $J_E$, and nonlinear frequency shift $\omega_1$ at the magnetic equator in the VHS code run. The dashed blue lines indicates the time when the formation of $-J_B$ begins, and the dashed red line indicates the time when $-J_E$ is formed, resulting in the nonlinear wave growth.
**Figure 5.** The optimum wave amplitudes (solid lines) with different values of $\tau$ (attached numbers) and the threshold for nonlinear wave growth (dashed line), (a) for simulation parameters used in Katoh and Omura [2007], and (b) for simulation parameters used in Hikishima et al. [2009].

**Figure 6.** (a) The optimum wave amplitudes for rising tone emissions (solid lines) with different values of $\tau$ (attached numbers), and the threshold of wave amplitude for the nonlinear wave growth (dashed line) and (b) the corresponding frequency sweep rates with the energetic electron density $N_h/N_c = 0.04$. Other physical parameters are specified for an observation by the Cluster spacecraft [Santolik et al., 2003; Santolik, 2008].
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