Parametric upconversion of lower hybrid wave by runaway electrons in tokamak

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A kinetic formalism of parametric decay of a large amplitude lower hybrid pump wave into runaway electron mode and an upper sideband mode is investigated. The pump and the sideband exert a ponderomotive force on runaway electrons, driving the runaway mode. The density perturbation associated with the latter beats with the oscillatory velocity due to the pump to produce the sideband. The finite parallel velocity spread of the runaway electrons turns the parametric instability into a stimulated Compton scattering process where growth rate scales as the square of the pump amplitude. The large phase velocity waves thus generated can potentially generate relativistic electrons. © *2010 American Institute of Physics*. [doi:10.1063/1.3442745]

I. INTRODUCTION

Relativistic runaway electrons (REs) are potentially a serious problem in advanced tokamaks such as International Thermonuclear Experimental Reactor (ITER).[1–4] REs with energies up to several MeV have been observed during numerous disruption in large tokamaks, such as Joint European Torus (JET),[5–8] Frascati Tokamak Upgrade (FTU),[9,10] Japan Torus (JT-60U),[11,12] and Hefei Tokamak-7 (HT-7).[13,14] They are a matter of serious concern as they can cause severe damage to the first wall structure on impact. REs can also arise in smaller numbers under normal tokamak operation, especially during startup and in low density plasmas. Two mechanisms are primarily attributed to runaway generation: primary generation through Dreicer accleration[15] and secondary generation provided by the avalanche effect.[16] For ITER disruptions, it has been predicted that avalanching would dominate and turn as much as two-thirds of the predisruption current into the runaway current.[17]

Current carrying fast electrons are generated/sustained by lower hybrid waves through parallel electron Landau damping when the Cerenkov resonance condition is fulfilled. Recently Martín-Solis *et al.*[9] observed experimentally large production of REs (up to ~80% of the predisruption plasma current) during a disruptive termination of discharge heated with lower-hybrid waves in FTU. Chen *et al.*[13] investigated the effect of lower hybrid waves on runaway production in the HT-7 tokamak, and showed that the presence of lower hybrid waves can greatly enhance the runaway production with high residual electric field. In Ref. 18, Liu and Mok proposed an elegant theory for the nonlinear evolution of RE distribution and time dependent synchrotron emission from tokamak.

In this paper we study the parametric upconversion of a lower hybrid wave into a low frequency electrostatic mode and a lower hybrid upper sideband. The low frequency mode is in resonant interaction with the REs. The lower hybrid pump wave $(\omega_0,\mathbf{k}_0)$ imparts an oscillatory velocity to electrons. When this is combined with a lower frequency density perturbation, a nonlinear current is produced driving the sideband $(\omega_2,\mathbf{k}_2)$. The sideband and the pump exert a ponderomotive force on electrons, driving the low frequency perturbation.

The paper is organized as follows. In Sec. II we obtain the RE susceptibility. Section III contains the nonlinear coupling, and growth rates have been calculated in seconds. Electron acceleration by the decay wave has been discussed in Sec. IV. Discussions have been given in Sec. V.

II. RE SUSCEPTIBILITY

We model the tokamak by a uniform plasma of background electron density n_0^0 in shearless magnetic field $\mathbf{B}=B_0\hat{z}$. The plasma has a component of REs of density n_{0r}^0. The initial RE distribution, before the wave-particle interaction, is determined from the kinetic equation with the boundary condition that satisfies the avalanche growth rate of the runaway density $dn_r/dt=n_r(E-1)/(c_z\tau_c)$,[16] giving

$$n_r = n_{r0}\exp[(E-1)t/(\tau_c c_z)], \tag{1}$$

where $E=eE_\parallel\tau_c/m_{e0}c$ is the normalized parallel electric field, assumed to be constant in time, c is the speed of light, $\tau_c=4\pi\epsilon_0^2 m_{e0}^2 c^3/n_e e^4 \ln\Lambda$ is the collision time for relativistic electrons, $c_z=\sqrt{3(Z+5)/\pi}\ln\Lambda$, Z is the effective ion charge, and n_{r0} is the seed produced by primary generation. Thus n_{0r}^0 is the value of runaway density n_r at the time when lower hybrid wave is launched. In a tokamak disruption, this initial distribution function of the relativistic tail of REs is[19,20]

$$f_0(p_z,p_\perp) = \frac{n_{r0}^0\alpha}{2\pi c_z p_z}\exp\left(-\frac{p_z}{c_z}-\frac{\alpha p_\perp^2}{2p_z}\right), \tag{2}$$

where $\mathbf{p}=\gamma\mathbf{v}/c$ is the normalized relativistic momentum of the REs and $\alpha=(E-1)/(Z+1)$, with γ as the relativistic factor.

We perturb this equilibrium by an elecrostatic perturbation

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$$\phi = A e^{-i(\omega t - \mathbf{k}\cdot\mathbf{r})}, \tag{3}$$

where $\omega \ll \omega_c$, $\omega_c = eB_0/m_{e0}$. The response of REs to it is governed by the Vlasov equation,

$$\frac{\partial}{\partial t} f + \mathbf{v}\cdot\nabla f = -\frac{e}{m_{e0}c}[\nabla\phi - \mathbf{v}\times\mathbf{B}_s]\cdot\nabla_{\mathbf{p}} f. \tag{4}$$

Writing $f = f_0 + f_1$ and linearizing Eq. (4) we obtain

$$f_1 = -\frac{e}{m_{e0}c}\int_{-\infty}^{t}\left[\nabla\phi\cdot\left(\frac{\partial}{\partial\mathbf{p}}f_0\right)\right]_{t'} dt', \tag{5}$$

where the integration is over the unperturbed trajectory. Equation (5) simplifies to give

$$f_1 = -\frac{e}{m_{e0}c}\phi\frac{n_{0r}^0\alpha}{2\pi c_z p_z}\exp\left(-\frac{p_z}{c_z} - \frac{\alpha p_\perp^2}{2p_z}\right) \times \sum_l J_l\left(\frac{k_\perp p_\perp c}{\omega_c}\right)\sum_n J_n\left(\frac{k_\perp p_\perp c}{\omega_c}\right) \times\left(\frac{k_z}{p_z} + \frac{k_z}{c_z} + \alpha\frac{l\omega_c}{c} - \alpha\frac{p_\perp^2}{2p_z^2}k_z\right)\frac{e^{i(l-n)\theta}}{\omega - \frac{k_z p_z c}{\gamma} - l\Omega}, \tag{6}$$

where J_l and J_n are the Bessel functions of order l and n, and θ is the gyrophase angle with $\Omega = eB/m_e = \omega_c/\gamma$. The perturbed density of REs turns out to be

$$n_b = \int_0^\infty\int_0^{2\pi}\int_{p_c}^\infty f_1 p_\perp dp_\perp d\theta dp_z, \tag{7}$$

where p_c is the boundary between the bulk and fast electron (tail region) momentum space.[21] For the anisotropy of the runaway distribution with relativistic electrons, with small argument of Bessel function, and by using the identity

$$\int_0^\infty e^{-s^2}J_l^2(\Psi s)s ds = \frac{1}{2}I_l e^{-\beta^2/2}, \tag{8}$$

the perturbed density takes the form

$$n_b = \frac{n_r e}{m_{e0}c_z}\frac{k_\perp^2}{\omega_c^2\alpha}\phi\gamma\left[\left(\frac{1}{c_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}\right)\frac{e^{-bp_c}}{b} + \left(\frac{D}{c_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}D - \frac{\omega_c\alpha}{k_zc}\right)e^{-bD}\lim_{\epsilon\to 0}\times\int_{(p_c-D)b}^\infty\frac{qe^{-q}}{q^2+b^2\epsilon^2}dq + i\pi e^{-bD}\times\left(\frac{D}{c_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}D - \frac{\omega_c\alpha}{k_zc}\right)\right], \tag{9}$$

where $b = 1/c_z + k_\perp^2c^2/\omega_c^2\alpha$, $D = (\gamma\omega - \omega_c)/k_zc$, and $q/b = p_z - D$.

For $(\omega \ll \omega_c)$ we can write

$$n_b = \frac{k^2}{e}\epsilon_0(\chi_{br} + i\chi_{bi})\phi, \tag{10}$$

where

$$\chi_{br} = \frac{\omega_{pr}^2}{c_z}\frac{k_\perp^2}{\omega_c^2\alpha k^2}\gamma\left[\left(\frac{1}{c_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}\right)\frac{e^{-bp_c}}{b} + \left(\frac{\gamma\omega}{k_zcc_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}\frac{\gamma\omega}{k_zc} - \frac{\omega_c\alpha}{k_zc}\right)e^{-b(\gamma\omega/k_zc)}\lim_{\epsilon\to 0}\times\int_{(p_c-\gamma\omega/k_zc)b}^\infty\frac{Qe^{-Q}}{Q^2+b^2\epsilon^2}dQ\right], \tag{11}$$

$$\chi_{bi} = \frac{\omega_{pr}^2}{c_z}\frac{k_\perp^2}{\omega_c^2\alpha k^2}\gamma\pi e^{-b(\gamma\omega/k_zc)}\left(\frac{\gamma\omega}{k_zcc_z} + \frac{k_\perp^2c^2}{2\omega_c^2\alpha}\frac{\gamma\omega}{k_zc} - \frac{\omega_c\alpha}{k_zc}\right),$$

where $\omega_{pr}^2 = n_{0r}^0e^2/m_{e0}\epsilon_0$ and $Q/b = p_z - \gamma\omega/k_zc$, $p_c = \sqrt{2T/m_{e0}}/c$.

III. NONLINEAR COUPLING AND GROWTH RATE

We consider the parametric coupling of a lower hybrid pump wave of potential

$$\phi_0 = A_0e^{-i(\omega_0t - \mathbf{k}_{0\perp}\cdot\mathbf{r} - k_{0z}z)}, \tag{12}$$

with a RE mode

$$\phi = Ae^{-i(\omega t - \mathbf{k}\cdot\mathbf{r})} \tag{13}$$

and an upper sideband mode

$$\phi_2 = A_2e^{-i(\omega_2t - \mathbf{k}_2\cdot\mathbf{r})}, \tag{14}$$

where $\omega_2 = \omega + \omega_0$ and $\mathbf{k}_2 = \mathbf{k} + \mathbf{k}_0$. The pump and sideband wave are primarily sustained by plasma electrons and ions and obey the dispersion relation

$$\omega_j^2 = \omega_{\mathrm{LH}}^2\left(1 + \frac{k_{zj}^2}{k_j^2}\frac{m_i}{m}\right), \tag{15}$$

where $j = 0, 2$, $\omega_{\mathrm{LH}}^2 = \omega_{pi}^2/(1 + \omega_p^2/\omega_c^2)$, $\omega_p = (n_0^0e^2/m_{e0}\epsilon_0)^{1/2}$, $\omega_{pi} = (Zn_0^0e^2/m_i\epsilon_0)^{1/2}$, and m_i is the ion mass. The ω, $\mathbf{k}$ mode has $\omega \simeq k_zv_{0b}^0$ and has prominent contribution from the REs. Here v_{0b}^0 is an average velocity of REs. We presume the pump to have parallel phase velocity less than v_{0b}^0. Let $\omega_0/k_{0z}v_{0b}^0 = \eta_0 < 1$, $\omega/k_zv_{0b}^0 = 1$, then $\omega_2/k_{2z}v_{0b}^0 = (\omega + \omega_0)/(k_z + k_{0z})v_{0b}^0 < 1$, i.e., the upper sideband also moves slower than the REs.

Had we considered lower hybrid pump wave to parallel phase velocity opposite to the velocity of the REs, $\omega_0/k_{0z}v_{0b}^0 = -\eta_0$, we would obtain

$$\omega_2/k_{2z}v_{0b}^0 = \frac{\frac{\omega}{\omega_0} + 1}{\frac{\omega}{\omega_0} - \frac{1}{\eta_0}} > 1, \tag{16}$$

i.e., the upper sideband moves faster than the runaways.

The pump and the sideband waves impart oscillatory velocity to plasma electrons

$$\mathbf{v}_{j\perp} = \frac{e}{m_{e0}\omega_c^2}[\boldsymbol{\omega}_c \times \nabla_\perp \phi_j - i\omega_j \nabla_\perp \phi_j],$$
$$v_{jz} = -\frac{ek_{jz}}{m_{e0}\omega_j}\phi_j, \tag{17}$$

where $j=0,2$.

The oscillatory velocity of REs can be obtained from the linearized equation of motion

$$m_{e0}\left[\frac{\partial}{\partial t}(\gamma \mathbf{v}_{bj}) + v_{0b}^0 \cdot \nabla(\gamma \mathbf{v}_{bj})\right] = e \nabla \phi_j - e\mathbf{v}_{bj} \times \mathbf{B}, \tag{18}$$

$$\mathbf{v}_{bj\perp} = \frac{e}{m_{e0}\omega_c^2}[\boldsymbol{\omega}_c \times \nabla_\perp \phi_j - i\gamma(\omega_j - k_{0z}v_{0bz}^0)\nabla_\perp \phi_j],$$
$$v_{bjz} = -\frac{ek_{jz}}{m_{e0}\gamma^3(\omega_j - k_{jz}v_{0bz}^0)}\phi_j. \tag{19}$$

The pump and upper sideband couples a low frequency ponderomotive force $\mathbf{F}_P$ on the electrons. $\mathbf{F}_P$ has two components: perpendicular and parallel to the magnetic field. The response of electrons to $\mathbf{F}_{P\perp}$ is strongly suppressed by the magnetic field and is usually weak. In the parallel direction, the electrons can effectively respond to $\mathbf{F}_{Pz}$; hence, frequency nonlinearity arises at $(\omega,\mathbf{k})$ mainly through $\mathbf{F}_{Pz}=-m\mathbf{v}\cdot\nabla v_z$. The parallel ponderomotive force, using the complex number identity $\mathrm{Re}\,\mathbf{A}\times\mathrm{Re}\,\mathbf{B}=(1/2)\mathrm{Re}[\mathbf{A}\times\mathbf{B}+\mathbf{A}^*\times\mathbf{B}]$, for the background electrons can be written as

$$\mathbf{F}_{pz} = eik_z\phi_p = -\left(\frac{m_{e0}}{2}\right)[\mathbf{v}_{0\perp}^* \cdot \nabla_\perp v_{2z} + \mathbf{v}_{2\perp} \cdot \nabla_\perp v_{0z}^*]. \tag{20}$$

Using Eq. (17) and considering only the dominant $\mathbf{E}_0\times\mathbf{B}$ drift terms, the ponderomotive potential ϕ_p takes the form

$$\phi_p = \frac{e\phi_0^*\phi_2}{2m_{e0}\omega_c^2}\frac{\mathbf{k}_{2\perp}\cdot\mathbf{k}_{0\perp}\times\omega_c}{ik_z\omega_2\omega_0}[\omega k_{0z} - \omega_0 k_z]. \tag{21}$$

The nonlinear density perturbation of the plasma electrons due to ponderomotive force can be written as

$$n^{NL} = -\frac{n_0^0 e k_z^2}{m\omega^2}\phi_p. \tag{22}$$

The linear density perturbation due to self-consistent potential ϕ is

$$n^L = (k^2\epsilon_0/e)\chi_e\phi,$$
$$\chi_e = \frac{\omega_p^2}{\omega_c^2}\frac{k_\perp^2}{k^2} - \frac{\omega_p^2}{\omega^2}\frac{k_z^2}{k^2}. \tag{23}$$

For the REs, the ponderomotive force can be written as

$$F_{pzb} = eik_z\phi_{pb} = -\left(\frac{m_{e0}\gamma}{2}\right)[\mathbf{v}_{b0\perp}^* \cdot \nabla_\perp v_{b2z} + \mathbf{v}_{b2\perp} \cdot \nabla_\perp v_{b0z}^*]. \tag{24}$$

Using Eq. (19) and considering only the dominant $\mathbf{E}_0\times\mathbf{B}$ drift terms, the ponderomotive potential ϕ_{pb} takes the form

$$\phi_{pb} = \frac{e\phi_0^*\phi_2}{2m_{e0}\omega_c^2\gamma^2}\frac{\mathbf{k}_{2\perp}\cdot\mathbf{k}_{0\perp}\times\boldsymbol{\omega}_c}{ik_z} \times\left(\frac{k_{0z}}{(\omega_0 - k_{0z}v_{0bz})} - \frac{k_{2z}}{(\omega_2 - k_{2z}v_{0bz})}\right). \tag{25}$$

One may note that the ponderomotive potential is maximum when $\mathbf{k}_\perp$ and $\mathbf{k}_{0\perp}$ are perpendicular to each other. The response of REs to the ponderomotive potential and the self-consistent potential ϕ

$$n_b = \frac{k^2}{e}\epsilon_0(\chi_{br} + i\chi_{bi})(\phi + \phi_{pb}). \tag{26}$$

Using Eqs. (22) and (26) in Poisson's equation $\nabla^2\phi=(e/\epsilon_0)\times(n+n_b-n_i)$, where $n=n^L+n^{\mathrm{NL}}$, we obtain

$$\epsilon\phi \cong -\chi_e\phi_p - \chi_b\phi_{pb}, \tag{27}$$

where $\varepsilon=1+\chi_e+\chi_b+\chi_i$.

The density perturbation at $(\omega,\mathbf{k})$ couples with the oscillatory velocity of electrons, $\mathbf{v}_0$, to produce nonlinear density perturbations at the upper sideband frequency. Solving the equation of continuity for the background electron,

$$\frac{\partial}{\partial t}n_2^{\mathrm{NL}} + \nabla\left(\frac{n}{2}v_0\right) = 0, \tag{28}$$

one obtains

$$n_2^{\mathrm{NL}} = \frac{n}{2\omega_2}(\mathbf{k}_2\cdot\mathbf{v}_0) \tag{29}$$

and for the REs

$$n_{2b}^{\mathrm{NL}} = \frac{n_b}{2\omega_2}(\mathbf{k}_2\cdot\mathbf{v}_{b0}). \tag{30}$$

Using Eqs. (29) and (30) in Poisson's equation for the upper sideband wave, we obtain

$$\varepsilon_2\phi_2 = \frac{k^2}{k_2^2}\frac{e\phi}{m_{e0}\omega_c^2}(1+\chi_i)\frac{\mathbf{k}_2\cdot\boldsymbol{\omega}_c\times\nabla_\perp\phi_0}{2\omega_2}, \tag{31}$$

where

$$\varepsilon_2 = 1 + \frac{\omega_p^2}{\omega_c^2} + \gamma\frac{\omega_{pr}^2}{\omega_c^2}\frac{(\omega_2 - k_{2z}v_{0bz})}{\omega_2} - \frac{\omega_{pi}^2}{\omega_2^2} \tag{32}$$

is the dielectric function at $(\omega_2,\mathbf{k}_2)$.

Equations (27) and (31) are the nonlinear coupled equations for ϕ and ϕ_2 from which nonlinear dispersion relation can be obtained

$$\epsilon\varepsilon_2 = \mu, \tag{33}$$

where the coupling coefficient can be written as

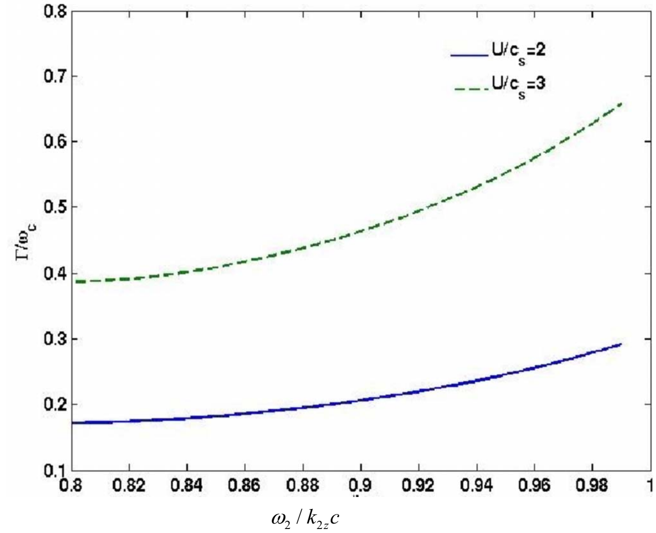

FIG. 1. (Color online) Variation in normalized growth rate as a function of normalized wavenumber for fast electron with energy of 100 keV.

$$\mu = \frac{U^2k^2\sin^2\delta}{4}\left[\frac{\chi_e}{\omega_2^2}\left(\frac{\omega k_{0z}}{\omega_0 k_z}-1\right) + \frac{\chi_b}{k_z\gamma^2\omega_2}\left(\frac{k_{0z}}{(\omega_0-k_{0z}v_{0bz})}-\frac{k_{2z}}{(\omega_2-k_{2z}v_{0bz})}\right)\right], \quad (34)$$

where $U=ek_0|\phi_0|/m_{e0}\omega_c$ is the magnitude of $\mathbf{E}_0\times\mathbf{B}$ electron velocity and δ is the angle between $\mathbf{k}_{2\perp}$ and $\mathbf{k}_{0\perp}$. We write $\omega=\omega_r+i\Gamma$, $\omega_2=\omega_{2r}+i\Gamma$, where ω_2 is the root of $\varepsilon_2=0$. Then Eq. (33) gives the growth rate

$$\Gamma = \frac{\mathrm{Im}[\mu(1+\chi_e+\chi_{br}+\chi_i-i\chi_{bi})]}{[(1+\chi_e+\chi_{br}+\chi_i)^2+\chi_{bi}^2]\dfrac{\partial\varepsilon_2}{\partial\omega_2}}. \quad (35)$$

In order to have a numerical appreciation of results, we consider the following set of parameters corresponding to HT-7 tokamak:[13] background electron density $\sim 4\times 10^{19}$ m^{-3}, temperature ~ 3 keV, ion temperature ~ 1.5 keV, magnetic field ~ 2.5 T, frequency of the lower hybrid pump ~ 2.45 GHz, and the density of the fast electron $\sim 2\times 10^{16}$ m^{-3}, $Z=1$, $\ln\Lambda=18$. In Fig. 1 we have plotted the normalized growth rate as function of normalized frequency by considering that the fast electron of energy ~ 100 keV, for different lower hybrid pump power $U/c_s=2$ and 3, where c_s is the ion sound speed, shows that the growth rate increases significantly with the increase in the lower hybrid power.

IV. ELECTRON ACCELERATION

The dynamics of a RE in the fast phase velocity lower hybrid sideband ϕ_2 can be described by the relativistic equation of motion

$$\frac{d\mathbf{p}}{dt} = -e[-\nabla\phi_2+\mathbf{v}\times\mathbf{B}_0]. \quad (36)$$

Expressing $d/dt=v_z d/dz$ the components of Eq. (36) can be written as

$$\frac{dp_x}{dz} = \frac{m_{e0}\gamma e k_\perp}{p_z}A_2\sin(\omega_2 t-k_{2x}x-k_{2z}z)-\omega_c m_{e0}\frac{p_y}{p_z},$$

$$\frac{dp_y}{dz} = m_{e0}\omega_c\frac{p_x}{p_z}, \quad (37)$$

$$\frac{dp_z}{dz} = \frac{m_{e0}\gamma}{p_z}ek_zA_2\sin(\omega_2 t-k_{2x}x-k_{2z}z).$$

These equations are supplemented with

$$\frac{dx}{dz} = \frac{p_x}{p_z},$$

$$\frac{dy}{dz} = \frac{p_y}{p_z}, \quad (38)$$

$$\frac{dt}{dz} = \frac{m_{e0}\gamma}{p_z}.$$

We introduce the dimensionless variables $P_x=p_x/m_{e0}c$, $P_y=p_y/m_{e0}c$, $P_z=p_z/m_{e0}c$, $X=x\omega_2/c$, $Y=y\omega_2/c$, $Z=z\omega_2/c$, $T=\omega_2 t$, and the electron drift velocity due to upper sideband $U_2=ek_{2\perp}A_2/m_{e0}\omega_c$; Eqs. (37) and (38) reduce to

$$\frac{dP_x}{dZ} = \frac{\omega_c U_2\gamma}{P_z c\omega_2}\sin\left(T-k_{2x}X\frac{c}{\omega_2}-k_{2z}Z\frac{c}{\omega_2}\right)-\frac{\omega_c P_y}{\omega_2 P_z},$$

$$\frac{dP_y}{dZ} = \frac{\omega_c P_y}{\omega_2 P_z},$$

$$\frac{dP_z}{dZ} = \frac{k_{2z}\omega_c U_2\gamma}{k_{2\perp}P_z c\omega_2}\sin\left(T-k_{2x}X\frac{c}{\omega_2}-k_{2z}Z\frac{c}{\omega_2}\right), \quad (39)$$

$$\frac{dX}{dZ} = \frac{P_x}{P_z},$$

$$\frac{dY}{dZ} = \frac{P_y}{P_z},$$

$$\frac{dT}{dZ} = \frac{\gamma}{P_z}.$$

We solve the above equations numerically for the following set of parameters: $P_x(0)=0.1$, $P_y(0)=0$, $P_z(0)=0.5$ $X(0)=0.5$, $Y(0)=0$, $Z(0)=0$, and normalized pump amplitude of the upper sideband wave $U_2/c_s=2$, $\omega_2/k_{2z}c=0.99c$, $k_{2z}/k_{2\perp}=1/25$. In Fig. 2 we have plotted the electron energy, normalized to the rest mass energy (on the gamma factor), as a function of distance of propagation Z. One will obtain large energy exchange between the particles and wave when phase synchronism condition is satisfied.

V. DISCUSSIONS

The free energy contained in the REs adds a quantum of energy to pump photon to produce the upper sideband photon. The REs open up the possibility of frequency upconversion of the lower hybrid pump wave. The parallel phase ve-

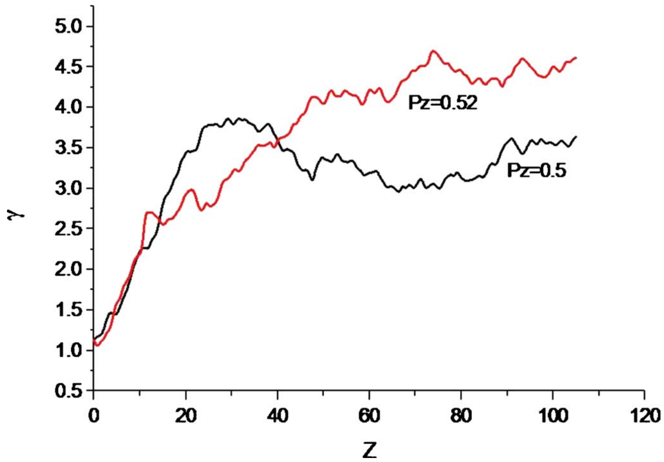

FIG. 2. (Color online) Variation in electron energy normalized to rest mass energy as a function of distance of propagation for $P_z \sim 0.5$ and 0.52.

locity of the upconverted lower hybrid wave is close to the velocity of light in vacuum. As the sideband wave acquires a large amplitude, comparable to that of the pump, it can accelerate the electrons to tens of MeV energy. For a typical case of electron drift velocity due to the upper sideband approaching twice the sound speed, the electron energy gain turns out to be ~ 2 MeV. The maximum growth rate occurs when δ is 90°, where δ is the angle between $k_{2\perp}$ and $k_{0\perp}$. However the other values of δ are possible, but they will give weaker growth rate.

The experiments on lower hybrid heating and current drive in tokamak have reported existence of MeV electrons that may be caused via parametric excitation of high parallel phase velocity waves. The energy gain by the electrons of $P_z \sim 0.5$ and 0.52 with the distance of propagation up to a point, after a while the particle is taken out of the resonance, by virtue of the energy gain, and it saturates. The energy gain is primarily through the Cerenkov resonance, though the transverse field also plays a role.

The energy gain is dependent to initial electron momentum. For given parameters, we obtain $P_z = 0.5 - 0.52$, where significant energy gain occurs. For 100 keV electrons, the Cerenkov resonance occurs when $\omega_2 = k_{2z} v_z$, and for $\omega_2 / k_{2z} = 0.99c$ the normalized energy gain $\gamma \sim 5$. This is the kind of energy gain we obtain in Fig. 2. In this paper we have ignored the toroidal and shear effects that may have profound effect on electron acceleration.

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