Role of Plasma Instabilities Driven by Oxygen Ions During Magnetic Storms and Substorms

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Ionospheric-origin O^+ ions consititute an important and some times dominant part of the ring current and the near-Earth plasma sheet region during geomagnetic storms and substorms. Low-frequency instabilities excited by the energetic oxygen ions in the near-Earth plasma sheet and ring current region during the geomagnetic storms/substorms are investigated. It is shown that the presence of ionospheric-origin oxygen ion beams with anisotropic pressure can excite helicon mode instability in the near-Earth plasma sheet region provided their Alfvénic Mach numbers lie in a certain range. The helicon modes are easily excited under the conditions when the usual long wavelengths fire-hose modes are stable. On the other hand, low-frequency quasi-electrostatic loss-cone type instabilities can be driven by the energetic oxygen ions in the storm time ring current region. These instabilities may scatter ring current particles and contribute to the ring-current decay.

INTRODUCTION

Observations in the various regions of the Earth's magnetosphere have shown the presence of protons and other energetic heavier ion distributions. Several observations suggest that ionospheric-origin O^+ ions consititute an important part of the inner magnetosphere and the near-Earth plasma sheet region [*Peterson et al.*, 1981; *Sharp et al.*, 1981; *Shelley et al.*, 1982; *Lockwood et al.*, 1985; *Delcourt et al.*, 1989; *Hultqvist*, 1991; *Daglis et al.*, 1991, 1993, 1994; *Cladis and Francis*, 1992; *Lennartsson*, 1994; *Wilken et al.*, 1995] during substorms/storms. Recently, it has been shown that the auroral ionospheric ion feeding of the inner plasma sheet during substorms can be fast (i.e., ~ characteristic substorm timescales), and that the ionosphere could actively influence the substorm energization processes by responding to the increased solar wind-magnetosphere coupling [*Daglis and Axford*, 1996]. The ionospheric O^+ ion flux in the near-Earth plasma sheet ($X \sim -6R_E$ to $-15\ R_E$) is observed to increase dramatically during the growth phase of substorms [*Daglis et al.*, 1990, 1991; *Kistler et al.*, 1992], and hence it could influence the dynamical evolution of the plasma sheet. However, the dominant mechanisms for the injection of ionospheric O^+ into the magnetosphere, and their subsequent acceleration to ring current energies during a storm are yet to be fully understood [*Daglis et al.*, 1999; *Hultqvist et al.*, 1999].

Disturbances in Geospace: The Storm-Substorm Relationship
Geophysical Monograph 142

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Lakhina and Tsurutani [1997, 1998] have shown that sufficiently strong fluxes of O^+ ions of ionospheric origin can excite an electromagnetic helicon mode instability in the near-Earth plasma sheet region. This instability can facilitate the excitation of tearing instability leading to substorm onset. During storms, plasma sheet oxygen ions can be accelerated by the helicon mode waves and injected earthwards and become part of the ring current.

It has been proposed that enhanced densities of ionospheric O^+ ions in some localized region in the plasma sheet may also lead to the excitation of the ion tearing instability [*Schindler*, 1974; *Baker et al.*, 1982], velocity shear instabilities [*Cladis and Francis*, 1992], or firehose type instabilities [*Verheest and Lakhina*, 1991; *Yoon et al.*, 1993; *Lakhina*, 1995, 1996], which could influence the substorm processes.

Recent observations by AMPTE/CCE have shown that energetic O^+ ions of ionospheric origin can become important constituents in the inner magnetosphere with the development of magnetic storms [*Krimigis et al.*, 1985; *Hamilton et al.*, 1988; *Roeder et al.*, 1996, *Daglis et al.*, 1993, *Daglis* 1997a,b]. During the geomagnetic quiet periods, the fractional concentration of the energetic oxygen ions in terms of energy density is about 10% [*Gloeckler and Hamilton*, 1987]. However, during the main phase of storms, the energy contents of the geomagnetically trapped particles increases and the fractional concentration of the oxygen ions is also increased. *Daglis* [1997a,b] has reported that the energy density of O^+ ions exceeds that of H^+ ions during intense storms with Dst of less than –200 nT, while the energy density of O^+ ions contributes only to 30% during a moderate magnetic storm. It is important to know how the thermal oxygen ions are extracted from the ionosphere and then energized to ring current energies of ~ a few keV to hundreds of keV. It has been suggested that the O^+ injection processes may involved either perpendicular heating or field-aligned acceleration with potential drops between the ionosphere and the magnetosphere [*Peterson et al.*, 1993].

An important issue related to the magnetic storm is the ring current decay. The ring current particles can be removed by various collisional processes as well as by scattering of charged particles by the waves. The loss mechanisms of ring current ions involving charge exchange with neutral atoms and Coulomb collisions with thermal plasma have been studied by several workers [*Fok et al.*, 1991, 1993, 1995; *Smith and Bewtra*, 1978; *Chen et al.*, 1997; *Ebihara and Ejiri*, 1999]. All ring current ions are subject to charge-exchange decay. However the decay rate depends on the ion mass and energy. It is found that high-energy O^+ are lost much faster than H^+, and field-aligned pitch-angle distributions experience larger losses than pancake like distributions [*Daglis et al.*, 1993]. The losses of heavier ions due to Coulomb collisions can become significant below energies of ~ 10 keV [*Fok et al.*, 1991; *Noel and Prolss*, 1997]. Large scale models incorporating the effects of charge exchange and Coulomb scattering tend to overestimate the flux of protons above tens of keV, and also yield a pitch-angle distribution which is too flat for energies of the order of 100 keV and higher [*Fok et al.*, 1995]. Pitch-angle scattering by the plasma waves has been suggested as a possible mechanism to account for these discrepancies [*Fok et al.*, 1996].

However, little work has been done on the role of wave-particle interaction on the ring current decay. Intense plasma waves could provide an efficient mechanism for energy transfer between different ion species and may prove important for selectively heating and accelerating thermal heavy ions [*Gendrin and Roux*, 1980; *Thorne and Horne*, 1994]. In fact, assessing the integrated effect on storm time ring current losses due to the scattering of ions by waves is one of the unsolved problem concerning the magnetic storm. The fact that the occurrence of particular plasma wave modes is usually limited in time or confined to localized regions, casts doubts on the ability of the wave scattering processes to significantly affect the energy balance of the ring current globally. The energetic protons and heavier ions can excite the low frequency waves such as electromagnetic ion cyclotron (EMIC) waves [*Cornwall et al.*, 1970, 1977; *Horne and Thorne*, 1993; *Kozyra et al.*, 1997]. Interaction between the ring current ions and the EMIC waves is considered as an important ring current loss process [*Cornwall et al.*, 1970, 1977; *Kozyra et al.*, 1997]. The time scale for scattering of ions into the loss cone during resonant interactions with EMIC waves can be rapid [*Lyons and Thorne*, 1972]. This could explain the rapid ring current energy loss on the time scales of 0.5–1.0 hrs during the main phases of intense-to-great geomagnetic storms [*Gonzalez et al.*, 1989; *Prigancova and Feldstein*, 1992].

In addition to the EMIC waves, quasi-electrostatic instabilities driven by loss-cone distributions of the ring current ions can also occur during magnetic storms [*Coroniti et al.*, 1972; *Bernstein et al.*, 1974; *Lakhina*, 1976, *Bhatia and Lakhina*, 1980a]. Recently, *Singh et al.* [2003] have studied quasi-electrostatic instabilities driven by anisotropic loss-cone oxygen ions. These instabilities can also scatter ring current ions rather efficiently leading to the ring current decay. These instabilities have not attracted much attention as compared to the EMIC waves, and there exist no global simulation model to ascertain the impact of these instabilities on the ring current energy budget.

HELICON MODES DRIVEN BY OXYGEN IONS IN THE NEAR-EARTH PLASMA SHEET REGION

In an electron-proton plasma, the dispersion relation for the right-hand polarized low-frequency modes, i.e., $\omega \ll \Omega_p$ (here ω and Ω_p represent the wave and the proton cyclotron frequencies, respectively), propagating parallel to the magnetic field, $\mathbf{B}_0$, gives the MHD Alfvén modes. In this case the proton Hall current completely cancels the electron Hall current, and the wave is maintained by the proton polarization current [*Papadopoulos et al.*, 1994]. However, in the presence of oxygen ions, the ion (both proton and oxygen) Hall currents cannot completely cancel the electron Hall current unless $\omega \ll \Omega_o$ (Ω_o being the oxygen ion cyclotron frequency). Therefore for the case when O^+ ions are weakly magnetized or unmagnetized, they carry negligible Hall current, and the resultant ion Hall current is not sufficient to neutralize the electron Hall current. This situation could give rise to helicon waves [*Papadopoulos et al.*, 1994; *Zhou et al.*, 1996; *Lakhina and Tsurutani*, 1997, 1998; *Lakhina*, 2001]. The helicon waves could lead to the fast current and flux penetration across the plasma sheet [*Papadopoulos et al.*, 1994], thus affecting substorm dynamics. *Lakhina and Tsurutani* [1997] showed that the presence of cold O^+ ion beams can excite helicon mode instability in the near-Earth plasma sheet region. *Lakhina* [2001] has shown that the helicon instability persists even when the O^+ ion beams are treated as hot.

Following *Lakhina* [2001], a dispersion relation for the electromagnetic modes propagating parallel to the magnetic field, $\mathbf{B_0} = B_0\mathbf{z}$ in a multispecies plasma can be written as,

$$\omega^2 = c^2k^2 - \sum_s \omega_{ps}^2\left[\frac{\omega - kU_s}{k\alpha_{\parallel s}}Z(\eta_s) + \left(\frac{\alpha_{\perp s}^2}{\alpha_{\parallel s}^2} - 1\right)\{1 + \eta_s Z(\eta_s)\}\right], \quad (1)$$

where $\omega_{ps} = (4\pi q_s^2 N_s/m_s)^{1/2}$ and $\Omega_s = q_s B_0/m_s c$ are the plasma and the gyrofrequency of the sth species, with $s = e$, p and o for electrons, protons and the oxygen ions respectively, N_s, q_s and m_s are the density, charge and mass of the sth species, U_s is the drift velocity of the sth species, c is the speed of light, and $\alpha_{\parallel s} = (2T_{\parallel s}/m_s)^{1/2}$, $\alpha_{\perp s} = (2T_{\perp s}/m_s)^{1/2}$ are respectively the parallel and perpendicular thermal velocities with respect to $\mathbf{B}_0$, $Z(\eta_s)$ is the plasma dispersion function with the argument $\eta_s = (\omega - kU_s \pm \Omega_s)/k\alpha_{\parallel s}$, and k is the wave number. The $\pm$ sign in η_s denotes the RH (+ sign) and the LH (– sign) modes. In writing (1), we have considered the drifting bi-Maxwellian distribution functions for the plasma species. In the equilibrium state, the charge neutrality is maintained by taking $N_e = N_p + N_o$.

In order to solve (1), we consider an unneutralized hot O^+ ion beam passing through the near-Earth plasma sheet region (corresponding to a finite field-aligned current in the system and $U_p = 0$, and $U_e = 0$). The background plasma is treated as cold, i.e., $\eta_e \gg 1$ and $\eta_p \gg 1$, and the ionospheric O^+ ion beam as hot, i.e., $\eta_o < 1$. Further, considering $(\omega - kU_s)^2 \ll \Omega_s^2$ for $s = e$ (electrons) and p (protons), and $\omega^2 \ll c^2k^2$, we obtain the following dispersion relation for the helicon modes

$$k^2V_{Ap}^2\left(1 - \frac{A_e}{2} - \frac{A_p}{2}\right) - R\Omega_o^2\left[-\frac{A_o}{\beta_{\parallel o}} + \frac{2\Omega_o(\Omega_o \mp kU_o)}{k^2\alpha_{\parallel o}^2}\right]$$
$$\mp R\Omega_o\omega\left(1 + \frac{2\Omega_o^2}{k^2\alpha_{\parallel o}^2}\right) - \omega^2 \quad (2)$$
$$-i\sqrt{\pi}R\Omega_o^2\exp(-\eta_o^2)\left[\frac{\beta_{\perp o}}{\beta_{\parallel o}}\frac{(\omega - kU_o)}{k\alpha_{\parallel o}} \mp \frac{A_o}{\beta_{\parallel o}}\frac{\Omega_o}{k\alpha_{\parallel o}}\right] = 0.$$

Here, $A_s = (\beta_{\parallel s} - \beta_{\perp s})$ denotes the pressure anisotropy of the plasma species and $\beta_{\parallel s}$ and $\beta_{\perp s}$ are respectively the parallel and the perpendicular plasma beta for the sth species, and $V_{Ap} = B_0/(4\pi\rho_p)^{1/2}$ is the Alfvén speed with respect to the proton mass density $\rho_p = N_p m_p$, and $R = (N_o m_o/N_p m_p)$ is the relative oxygen ion mass density with respect to protons.

Assuming $\omega^2 \ll R\Omega_o\omega$, and writing $\omega \ll \omega_r + i\gamma$, where $\gamma^2 \ll \omega_r^2$, then from (2), we obtain the following expressions for the real frequency, ω_r, and growth rate, γ,

$$\omega_r = \pm\frac{\left\{\frac{N_e}{N_o}\left(1 - \frac{A_e}{2} - \frac{A_p}{2}\right)\frac{k^2c^2}{\omega_{pe}^2}\Omega_e + \Omega_o\left[\frac{A_o}{\beta_{\parallel o}} - \frac{2\Omega_o(\Omega_o \mp kU_o)}{k^2\alpha_{\parallel o}^2}\right]\right\}}{\left(1 + \frac{2\Omega_o^2}{k^2\alpha_{\parallel o}^2}\right)}, \quad (3)$$

$$\gamma \approx \frac{\sqrt{\pi}\Omega_o\exp(-\eta_{or}^2)}{\left(1 + \frac{2\Omega_o^2}{k^2\alpha_{\parallel o}^2}\right)}\left[\frac{\beta_{\perp o}}{\beta_{\parallel o}}\frac{(\omega_r - kU_o)}{k\alpha_{\parallel o}} \mp \frac{A_o}{\beta_{\parallel o}}\frac{\Omega_o}{k\alpha_{\parallel o}}\right], \quad (4)$$

where $\eta_{or} = (\omega_r - kU_o \pm \Omega_o)/k\alpha_{\parallel o}$. The $\pm$ sign in (3) and η_{or} denotes the RH (+ sign) and the LH (– sign) modes.

When the second term in the curly brackets is neglected, (3) becomes very similar in structure to the dispersion relation for the usual helicon mode,

$$\omega_H = \frac{k^2c^2}{\omega_{pe}^2}\Omega_e, \quad (5)$$

in electron-proton plasma. Equation (5) shows that electron dynamics dominates the interaction between the electromagnetic waves and the multi-ion plasma.

Figure 1 shows the real frequency, ω_r, and the growth rate, γ versus the normalized wave number for the helicon modes for some typical parameters of the near-Earth plasma sheet. For parameters listed in the caption of Figure 1, it is found that the helicon mode instability occurred for the RH mode only. The assumption of treating the O^+ ion beam as hot breaks down at smaller values of normalized wavenum-

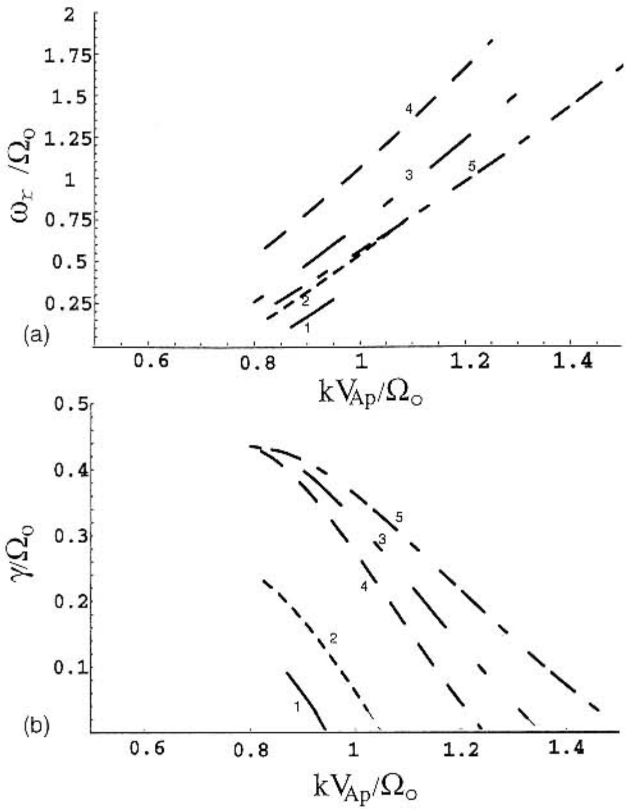

Figure 1. Variation of normalized real frequency ω_r/Ω_o (a), and growth rate γ/Ω_o (b) versus normalized wavenumber kV_{Ap}/Ω_o for the helicon mode instability driven by non-neutralized O^+ ions in the CPS region for $M = U_o/V_{Ap} = 0.25$, R=1.0, $A_e = A_p$=0, and $\beta_{\|o} = 3.5$. The curves 1, 2, and 3 are respectively for $A_o = (\beta_{\|o} - \beta_{\perp o})$ = 0.1, 0.5, and 2.0. The curves 4 and 5 are for A_o = 2, M=0.25, R=1.0, and A_p=−0.5 and 0.5, respectively. For the parameters considered here, the LH mode instability does not exist.

bers, kV_{Ap}/Ω_o where the peak in the growth rate occurs. Therefore in Figure 1, only those portion of the growth rate (and the corresponding real frequency) curves where all the assumption made in the analysis are satisfied are shown. It is clear from Figure 1 that ranges of real frequencies and growth rate increase when A_o is increased (cf. Curves 1, 2 and 3). Positive (negative) value of proton anisotropy has a destabilizing (stabilizing) effect (cf. Curves 4 and 5). Similar analysis for the case of neutralized O^+ ion beam (i.e., no field-aligned current) shows that in general, both the growth rate and real frequency are smaller as compared to the non-neutralized O^+ ion beam case. Further, the instability occurs for higher values of the normalized wave number in the case of neutralized O^+ ion beam as compared to the non-neutralized O^+ ion beam case. Therefore, the unstable modes for the neutralized beam case would have shorter wavelengths as compared to the latter case.

We consider the following parameters for the central plasma sheet (CPS) in the region $X \sim -10\ R_E$ to $-15\ R_E$: A_0 =0.1 to 2.0 [*Daglis et al.*, 1991; *Lennartsson*, 1994; *Cladis and Francis*, 1992], R = 1 – 10 [*Wilken et al.*, 1995], and $|U_o - U_p| \approx 10 - 60$ km s^{-1} [*Peterson et al.*, 1981; *Orsini et al.*, 1985]. Then, for B_0 = 10 nT and N_p = 0.5 cm^{-3}, we get typical Mach numbers $M = U_o/V_{Ap}$ = 0.025–0.25 in the CPS region. Then, the real frequencies, growth rates, and unstable wavelengths, $\lambda = 2\pi/k$, for the helicon mode instability are found to be ω_r = (1.0 – 19.0) mHz, γ = (0.5 – 4.5) mHz, and λ = (0.5 – 2) R_E, respectively. The typical *e*-folding time of the instability is about 3 to 15 minutes at wavelengths of $\lambda \approx 0.5$ to 2 R_E, which is reasonably short. Therefore, these modes could attain saturation as the enhanced convection events may last for a few hours.

The large amplitude helicon mode may be responsible for some of the low-frequency RH polarized electromagnetic noise in the ULF – ELF frequency range observed in the CPS and plasma sheet boundary layer [*Russell*, 1972; *Tsurutani et al.*, 1985, 1987; *Bauer et al.*, 1995). The large scale fluctuating z and y components of B, associated with the helicon modes could twist the equilibrium magnetic field into flux ropes. Further, the large amplitude helicon modes could produce localised minima in B near the neutral axis that are the potential site for the excitation of the tearing mode instabilities. The tearing instabilities could lead to the onset of the expansion phase of substorms. Moreover, the low-frequency turbulence due to the helicon modes could scatter electrons trapped in the CPS region, thus further enhancing the growth of the tearing modes at these sites.

It has recently been noted that O^+ and H^+ ions in the dayside boundary layer/cusp are heated/accelerated essentially all of the time by intense broadband plasma waves. During southward interplanetary magnetic field events, these ions could be convected in the anti-sunward direction across the polar cap and into the plasma sheet [*Tsurutani et al.*, 1998]. There can also be a direct injection of O^+ ions from the auroral region as suggested by [*Daglis et al.*, 1994]. Such energetic oxygen ions could provide an ideal situation for the excitation of the helicon mode instability and substorm onset. These O^+ ions would be energized, along with the H^+ and other minor ions, during the substorm expansion phase. Consequential injection of this plasma inward into the Earth's night side magnetosphere by the storm electric fields could then lead to enhanced oxygen ion fluxes in the storm-time ring current. The enhancement of the O^+ ions in the ring current region could have a strong controlling effect on the excitation of EMIC waves, and thus on the ring current decay during storms [*Thorne and Horne*, 1997].

RING CURRENT INSTABILITIES DRIVEN BY ENERGETIC OXYGEN IONS

Electromagnetic Ion Cyclotron Waves

The AMPTE/CCE spacecraft frequently observed EMIC waves in the outer magnetosphere beyond $L = 7$ [*Anderson et al.*, 1990, 1992]. The wave spectral properties and frequency of occurrence depends on local time and can be explaind by the resonant instability of anisotropic ring current H^+ ions in the variable background thermal plasma [*Cornwall*, 1970; *Horne and Thorne*, 1993, 1994). The presence of oxygen ions during the storms can affect the excitation of EMIC and other low-frequency instabilities. *Thorne and Horne* [1994] showed that during the geomagnetic storms when the oxygen ion content is significantly enhanced, the absorption of EMIC waves become efficient and may lead to the acceleration of O^+ ions of ionospheric origin to ring current energies. *Horne and Thorne* [1994] studied the convective instabilities of the EMIC waves in the outer magnetosphere and discussed that the waves below the cyclotron frequency of the helium (He^+) ions are not reflected when the O^+ ion concentration is small. *Thorne and Horne* [1997] studied the modulation of EMIC instability due to the interaction with ring current O^+ ions during a geomagnetic storms. They demonstrated that during the modest storm, strong EMIC excitation can occur in the frequency band above the oxygen gyrofrequency due to cyclotron resonance with the anisotropic ring current H^+ ions. The energy of the excited wave is efficiently converted into the perpendicular heating of the O^+ ions thus making it more anisotropic during the main phase of the magnetic storm.

Jordanova et al. [1996] have modeled the global impact of ion cyclotron waves on the ring current by using the Ring Current-Atmosphere Interaction Model (RAM) which follows the evolution of three major ring current ion species (H^+, He^+, and O^+) considering adiabatic drift motion, Coulomb collisions, charge exchange, and pitch-angle scattering of protons in the field of EMIC waves. The model produced order- of -magnitude enhancements in the ion precipitation as a result of diffusion in the ion cyclotron waves within the unstable region. However, no significant impact of the wave losses was seen in the global energy balance though the waves reduced the anisotropy in the proton pitch angle distributions locally. *Kozyra et al.* [1997] implemented an improved scheme for estimating the global distribution and amplitude of EMIC activity using a warm plasma ray tracing program (HOTRAY) [*Horne and Thorne*, 1993] in the RAM model. The integrated energy loss from the ring current, due to the scattering of protons into the loss cone, caused an additional ~ 8 nT recovery in the Dst index and thus was found to be important for the global energy balance of the ring current.

Low-Frequency Quasi-Electrostatic Instabilities

Electrostatic noise in the frequency range of (20–500) Hz has been reported by *Anderson and Gurnett* [1973] just outside the plasmapause. *Labelle et al.* [1988] have observed electrostatic noise from (30–100) Hz in the ring current region. *Olsen et al.* [1987] observed a broadband electrostatic noise in frequency range of 1.8 Hz to 500 Hz and beyond peaking near the lower hybrid frequency (~ 100 Hz) in the ring current region. A quasi-electrostatic instability excited by loss-cone distribution of ring current protons [*Coroniti et al.*, 1972, *Bernstein et al.*, 1974, *Lakhina*, 1976, *Bhatia and Lakhina*, 1980a] can generate frequencies close to the lower hybrid frequency or below the ion plasma frequency. The instability is found to be strongest for $\beta \sim 1$, low-density region just outside the plasmapause and is nearly stable for auroral regions and high density plasmasphere for $\beta >> 1$ [*Coronoti et al.*, 1972].

We study the obliquely propagating low-frequency quasi-electrostatic waves generated by the anisotropic oxygen ions in the ring current region during the main phase of the storms. Ring current plasma is considered to be consisting of Maxwellian distributed electrons, and protons, and energetic oxygen ions having Dory-Guest-Harris (DGH) [*Dory et al.*, 1965, *Coroniti et al.*, 1972] type of distribution given by

$$f_o = \frac{N_o}{\pi^{3/2}} \frac{1}{J!\,\alpha_{\|o}^3} \left(\frac{\upsilon_\perp}{\alpha_{\|o}}\right)^{2J} exp\left[-\frac{\upsilon_\perp^2 + \upsilon_\|^2}{\alpha_{\|o}^2}\right], \qquad (6)$$

where $J = \left(\frac{T_{\perp o}}{T_{\|o}} - 1\right)$ defines the temperature anisotropy, and $\alpha_{\|o} = (2T_{\|o}/m_o)^{1/2}$ is the parallel thermal velocity of the oxygen ions with respect to ambient magnetic field $\mathbf{B}_o = B_o\mathbf{z}$. Here, we treat electrons and protons as magnetized, and oxygen ions as unmagnetized, i.e., $\Omega_o < \omega < \Omega_p << \Omega_e$. The protons and oxygen ions are considered to be electrostatic, i.e., $\omega_{pp}^2/c^2k^2 << 1$, and $\omega_{po}^2/c^2k^2 << 1$ and electrons response is electromagnetic, i.e., $\omega_{pe}^2/c^2k^2 >> 1$. Under the above approximations, a linear dispersion relation for the low-frequency waves propagating obliquely to the magnetic field, $\mathbf{B}_o$ can be obtained by solving the linearized Vlasov equation and Maxwell's equations and is written as (*Davidson et al.*, 1977, *Bhatia and Lakhina*, 1977, 1980a, b)

$$1+\frac{\omega_{pe}^2}{\Omega_e^2}\left[\frac{1-I_0(\lambda_e)e^{-\lambda_e}}{\lambda_e}\right]+\frac{\omega_{pe}^2}{\Omega_e^2}\frac{\omega_{pe}^2}{c^2k^2}\left[\frac{[\{I_0(\lambda_e)-I_1(\lambda_e)\}e^{-\lambda_e}]^2}{1+\beta_e\{I_0(\lambda_e)-I_1(\lambda_e)\}e^{-\lambda_e}}\right]$$
$$-\frac{k_\parallel^2}{k_\perp^2}\frac{\omega_{pe}^2}{\omega^2}\frac{I_0(\lambda_e)e^{-\lambda_e}}{1+\frac{\omega_{pe}^2}{c^2k^2}I_0(\lambda_e)e^{-\lambda_e}}-\sum_{n=1}^{\infty}\frac{4n^2\omega_{pp}^2\Omega_p^2}{k_\perp^2\alpha_p^2(\omega^2-n^2\Omega_p^2)}I_n(\lambda_p)e^{-\lambda_p} \quad (7)$$
$$-\frac{\omega_{po}^2}{k^2\alpha_{\parallel o}^2}\frac{(-1)^J}{J!}\mu^J\left[\frac{d^J}{d\mu^J}[Z'\left(\frac{\omega}{k}\sqrt{\mu}\right)]\right]=0,$$

where $Z'\left(\frac{\omega}{k}\sqrt{\mu}\right)$ is the derivative of the plasma dispersion function with respect to its argument, and $\lambda_{e,p} = k_\perp^2\alpha_{e,p}^2/2\Omega_{e,p}^2$, $\mu^{-1} = \alpha_{\parallel o}^2$ and I_n is the modified Bessel function of order n, and $\alpha_{e,p} = (2T_{e,p}/m_{e,p})^{1/2}$ is the thermal velocity of electrons and protons respectively. We expand the plasma dispersion function in the limit $\omega\sqrt{\mu}/k \leq 1$ for the resonant oxygen ions, and consider $n = 1$ terms for the protons in the above dispersion relation (7). The latter assumption restricts the analysis to frequencies smaller than the proton cyclotron frequencies. Substituting for $\omega = \omega_r + i\gamma$ in (7) and separating real and imaginary parts, we obtain following expressions for real frequency (ω_r) and the growth rate (γ),

$$\omega_r^2 = \frac{B\pm\sqrt{B^2-4C}}{2}, \quad (8)$$

and,

$$\gamma = -\frac{\sqrt{\pi}}{P}\frac{\omega_r^2\omega_{po}^2}{\alpha_{\parallel o}^3k^3}\left[\frac{(-1)^J}{J!}\mu^{J-1/2}\frac{d^J}{d\mu^J}\left\{\mu^{1/2}e^{-\mu\frac{\omega_r^2}{k^2}}\right\}\right], \quad (9)$$

where,

$$B = \left[\Omega_p^2+\frac{\omega_{pe}^2}{P}\frac{k_\parallel^2}{k_\perp^2}\frac{I_0(\lambda_e)e^{-\lambda_e}}{1+\frac{\omega_{pe}^2}{c^2k^2}I_0(\lambda_e)e^{-\lambda_e}}+\frac{4}{P}\frac{\Omega_p^2\omega_{pp}^2}{k_\perp^2\alpha_p^2}I_1(\lambda_p)e^{-\lambda_p}\right], \quad (10)$$

$$C = \frac{1}{P}\frac{k_\parallel^2}{k^2}\omega_{pe}^2\Omega_p^2\frac{I_0(\lambda_e)e^{-\lambda_e}}{1+\frac{\omega_{pe}^2}{c^2k^2}I_0(\lambda_e)e^{-\lambda_e}}, \quad (11)$$

and,

$$P \simeq 1+\frac{\omega_{pe}^2}{\Omega_e^2}\left[\frac{1-I_0(\lambda_e)e^{-\lambda_e}}{\lambda_e}\right]+\frac{\omega_{pe}^2}{\Omega_e^2}\frac{\omega_{pe}^2}{c^2k^2}\left[\frac{[\{I_0(\lambda_e)-I_1(\lambda_e)\}e^{-\lambda_e}]^2}{1+\beta_e\{I_0(\lambda_e)-I_1(\lambda_e)\}e^{-\lambda_e}}\right]$$
$$+2\frac{\omega_{po}^2}{k^2\alpha_{\parallel o}^2}\frac{(-1)^J}{J!}\mu^J\sigma \quad (12)$$

where $\sigma = 1$ for $J = 0$ and $\sigma = 0$ for $J \geq 1$. It may be noticed from the growth rate expression (9) that for $J = 0$, i.e., when perpendicular and parallel temperatures of oxygen ions are equal, the γ is negative and hence, waves are damped and the system is stable. It is expected as for $J = 0$, the plasma system is isotropic and there is no free energy source to excite the instability. For $J \geq 1$, the energetic oxygen ions have finite anisotropy which can drive the low frequency instability.

The growth rate expression for $J = 1$ is given by

$$\gamma = \frac{\sqrt{\pi}}{2P}\frac{\omega_r^2\omega_{po}^2}{\alpha_{\parallel o}^3k^3}\left[1-2\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}\right]e^{-\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}}, \quad (13)$$

and for $J = 2$ growth rate is given by

$$\gamma = \frac{\sqrt{\pi}}{8P}\frac{\omega_r^2\omega_{po}^2}{\alpha_{\parallel o}^3k^3}\left[1+4\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}-4\frac{\omega_r^4}{k^4\alpha_{\parallel o}^4}\right]e^{-\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}}, \quad (14)$$

and similarly for $J = 3$ growth rate is given by

$$\gamma = \frac{\sqrt{\pi}}{16P}\frac{\omega_r^2\omega_{po}^2}{\alpha_{\parallel o}^3k^3}\left[1+2\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}+4\frac{\omega_r^4}{k^4\alpha_{\parallel o}^4}-\frac{8}{3}\frac{\omega_r^6}{k^6\alpha_{\parallel o}^6}\right]e^{-\frac{\omega_r^2}{k^2\alpha_{\parallel o}^2}}, \quad (15)$$

It can be seen from the growth rate expression (13) for $J = 1$, that for the growth of the waves a condition $\omega_r^2/k^2\alpha_{\parallel o}^2 < 1/2$ must be satisfied which is consistent with our assumption used in expanding the plasma dispersion function for the oxygen ions. The chosen parameters for numerical calculations for $J = 1$ case, are $k_\parallel/k = 0.05$, $T_{\parallel o}/T_e = 30$, $T_{\parallel o}/T_p = 1$, $\omega_{pe}/\Omega_e = 2$, $T_e = 1$ keV, $N_e = 10$ cm^{-3} and $\alpha_{\parallel o}/c \simeq 2 \times 10^{-3}$ which are representative of storm time ring current, and have been used for the results presented in the Figures 2–4.

Figure 2 shows the variation of normalized real frequency ω_r/Ω_o with $\lambda_o = k_\perp^2\alpha_{\parallel o}^2/2\Omega_o^2$ for fractional oxygen ion concentration, $N_o/N_e = 0.2$. The real frequency shows marginal increase with the increase of oxygen ion concentration. Figure 3 shows the variation of normalized growth rate γ/Ω_o with λ_o for various values of the fractional oxygen ion density N_o/N_e as shown on the curves. It may be noted that as the oxygen ion concentration increases, the growth rate increases. The range of unstable wavenumbers decreases with increase in oxygen ion concentration.

Figure 4 shows the comparison of the growth rates for $J = 1$, 2 and 3 values of the temperature anisotropies for the above mentioned parameters. The peak growth rate increas-

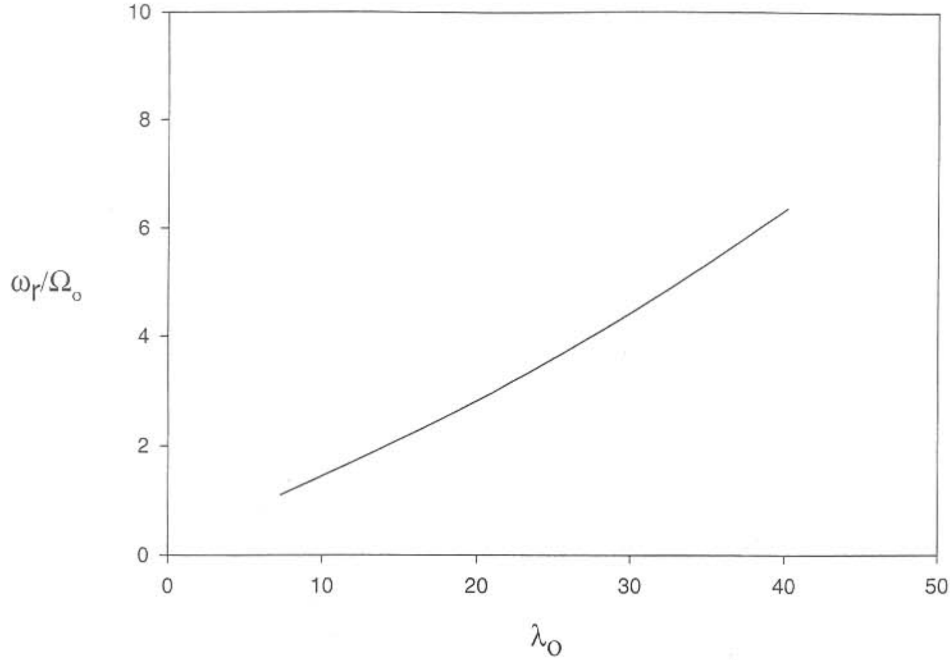

Figure 2. shows the normalized real frequency ω_r/Ω_o versus $\lambda_o = k_\perp^2 \alpha_{\parallel o}^2/2\Omega_o^2$ for $J = 1$ for fractional oxygen ion concentration $N_o/N_e = 0.2$. The typical parameters are, $k_\parallel/k = 0.05$, $T_{\parallel o}/T_e = 30$, $T_{\parallel o}/T_p = 1$, $\omega_{pe}/\Omega_e = 2$, and $\alpha_{\parallel o}/c \simeq 2 \times 10^{-3}$ for $J = 1$ case.

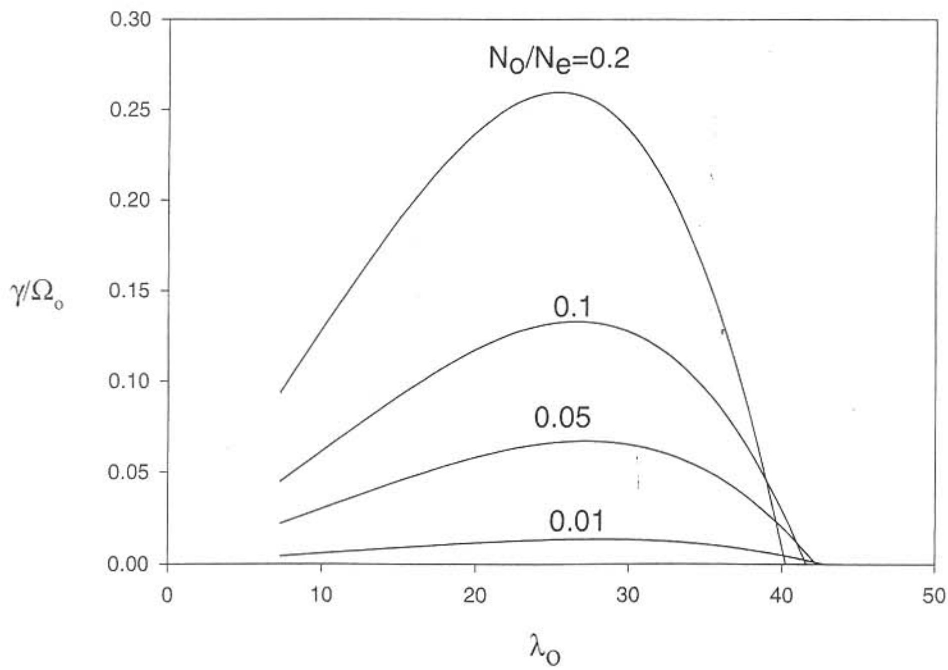

Figure 3 shows the variation of normalized growth rate γ/Ω_o with λ_o for various values of the fractional oxygen ion density N_o/N_e as shown on the curves. Other parameters are same as in Figure 2.

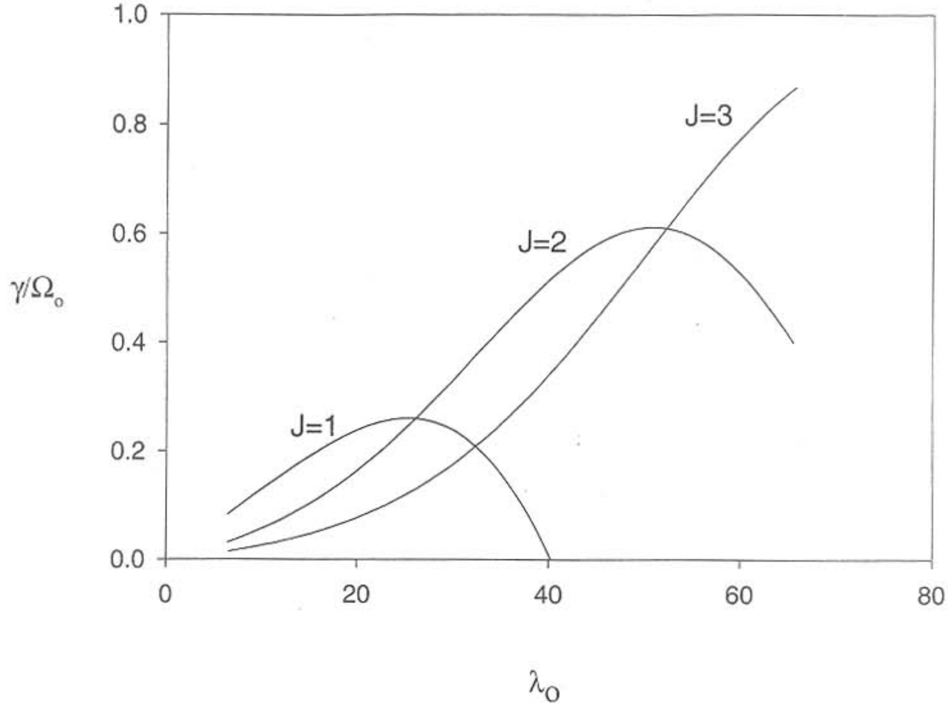

Figure 4 shows the comparison of the growth rates for $J = 1$, 2 and 3 for $N_o/N_e = 0.2$. Other parameters are same as for Figure 2.

es with the increase of the anisotropy. However, real frequency is not affected by the change in anisotropy. It is also found that the range of unstable wavenumbers increases with the increase in the anisotropy.

Thorne and Horne [1997] have provided the Maxwellian fits to the energetic H^+ and O^+ ion distributions obtained from the ring-current atmosphere interaction model (RAM) and have given the equatorial parameters. We have carried out the numerical calculations of the growth rate, real frequency etc. for the three different values of anisotropic index J for the storm time ring-current parameters. For $J =$ 1, the frequency of the wave $\omega_r \sim (0.5 - 3)$ Hz and the corresponding to the growth rate $\gamma \sim (40 - 125)$ mHz as shown in Figure 4. The corresponding perpendicular wavelength $(\lambda_\perp)$ of the excited waves is $\sim (140 - 350)$ km. For $J = 2$, the corresponding frequency and growth rate of the waves are in the range $\sim (0.5 - 5.5)$ Hz and $\sim (15 - 290)$ mHz respectively. The perpendicular wavelength of these waves is $\sim$ (109 – 350) km. For $J \approx 3$, the real frequency and growth rate are $\sim (0.5 - 5.5)$ Hz and $\sim (7 - 400)$ mHz respectively for $\lambda_\perp \sim (109 - 350)$ km. Thus, it can be concluded that for the higher values of anisotropy, the peak growth rate will be greater.

The saturation electric field amplitude of the waves can be estimated by comparing the linear wave growth rate with the trapping frequency of the oxygen ions [*Coroniti et al.*, 1972]. Thus, assuming that the modes are stabilized by the trapping of O^+ ions by the waves when their amplitude becomes large, the saturation electric field (E_s) for the low-frequency waves is given by

$$E_s = \frac{\pi m_o \omega_r^4 \omega_{po}^4}{e \alpha_{\parallel o}^4 k^7 P^2} \left[\frac{(-1)^J \mu^J}{J!} \frac{d^J}{d\mu^J} \left\{ \sqrt{\mu} \exp\left(-\frac{\omega_r^2}{k^2} \mu \right) \right\} \right]^2 . \quad (16)$$

To estimate the diffusion time for an oxygen ion to reach the loss cone, a diffusion coefficient $D = (\Delta v_\perp)^2/2\Delta t$ is constructed, where $\Delta v_\perp = eE_s \Delta t/m_o$ and Δt is the wave particle correlation time. If we assume $\Delta t \approx \gamma^{-1}$, then D is given by

$$D = \frac{\gamma^3}{2k^2} \quad (17)$$

The saturation electric fields for the $J = 1$, 2, and 3 as calculated from (16) are $\sim (0.5 - 3)$ mV/m, $\sim (80\mu\text{V/m} - 11$ mV/m) and $\sim (0.2\ \mu\text{V/m} - 20$ mV/m) respectively and the corresponding diffusion coefficients as calculated from (17) are of the order of $\sim (2.4 \times 10^{11} - 2 \times 10^{12})$ cm²/s³, $\sim (1.3 \times 10^{10} - 1.2 \times 10^{13})$ cm²/s³ and $\sim (1.25 \times 10^9 - 2.65 \times 10^{13})$ cm²/s³ respectively. Thus, saturation electric field and diffusion coefficients will increase with the increase in the oxygen ion anisotropy. The time during which an oxygen ion can diffuse into loss cone with a velocity $\alpha_{\|o}$ is ~ (10^2 – 10^6) s.

CONCLUSION

The ionospheric origin oxygen ions dominate the ring current and near Earth plasma sheet region during the storms and substorms. The oxygen ion beams with anisotropic pressure can excite the helcon mode instability for the RH mode in the near Earth plasma sheet region. The e-folding time of the instability is about (3–15) minutes and these modes could attain saturation as the enhanced convection events last few hours. The low-frequency turbulence due to the helicon modes could scatter electrons trapped in the central plasma sheet region and enhance the growth of the tearing modes.

Presence of O^+ ions during the storms facilitates the absorption of EMIC waves and lead to the acceleration of O^+ ions of ionospheric origin to ring current energies. During the storms the EMIC waves can heat the O^+ ions and thus make them more anisotropic.

It is possible to excite the low-frequency quasi-electrostatic wave with anisotropic oxygen ions. The frequency of these waves is found to be between the oxygen ion and proton cyclotron frequencies. These waves have larger growth rate for nearly perpendicular propagation. Low-frequency waves are found to have a higher growth rate in the presence of higher anisotropy of oxygen ions but the corresponding real frequency is not affected by the anisotropy. These quasi-electrostatic low frequency waves considered here may explain the lowest frequency component (< 10 Hz) of the electrostatic waves observed in the ring current region beyond the plasmapause by *Olsen et al.* [1987]. Growth time of these waves is few tens of seconds whereas geomagnetic storm main phase lasts for several hours. The low-frequency waves may scatter the ring-current particles and lead to the ring-current decay. Thus, the interaction between the particles and the low-frequency waves studied here can provide a mechanism for the ring current decay that is complimentary to the charge exchange mechanism. However, the study has not taken into account the effect of cold plasma which become necessary when hot ring current penetrate the plasmapause during the main phase of the magnetic storm. We expect that the cold plasma population may reduce the growth rate of the low frequency waves in the ring current region as the fractional concentration of oxygen ions will be reduced. There is also a possibility that the presence of cold plasma species may allow the excitation of new modes. However, a detailed study is required to find out the effect of cold plasma on the growth rate and the range of the unstable wave numbers of the low frequency quasi-electrostatic wave modes. Further, our analysis treats the ring current system as a uniform plasma. This assumption is justified since the typical transverse wavelengths, $\lambda_\perp$ ~ 100–350 kms are much smaller than the inhomogeneity scale lengths in the ring current region.

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