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# On the onset of HF-induced airglow at HAARP



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[1] Observations of airglow at 630 nm (red line) and 557.7 nm (green line) during the February 2002 campaign at the High Frequency Active Auroral Research Program (HAARP) heating facility are analyzed. We find that during injections toward magnetic zenith (MZ) the green and red lines gain ~5 R within ~1 s and ~20 R within ~10 s, respectively. We term this period the onset of the HF-induced airglow. A model of the onset at magnetic zenith is developed. It accounts for background photoelectrons and dissociative recombination of $O_2^+$. It is shown that heating and acceleration of background electrons dominate the airglow onset. We propose a scenario for the generation of strong Langmuir turbulence for injections outside the Spitze region, including magnetic zenith. *INDEX TERMS:* 2403 Ionosphere: Active experiments; 2483 Ionosphere: Wave/particle interactions; 2471 Ionosphere: Plasma waves and instabilities; 2481 Ionosphere: Topside ionosphere; *KEYWORDS:* HF modification experiments, artificial airglow onset

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## 1. Introduction

[2] A distinctive feature of HF modification experiments is the excitation of airglow at 630.0 and 557.7 nm by high-power, high-frequency (HF) radio waves [e.g., *Sipler et al.*, 1974; *Bernhardt et al.*, 1989; *Pedersen and Carlson*, 2001; *Gustavsson et al.*, 2001, 2002; *Kosch et al.*, 2000, 2002]. Enhancements up to ~500 R (Rayleighs) with the green-to-red ratio $\zeta_{gr}$ as high as $\geq 0.3$ have been reported [*Gustavsson et al.*, 2002; *Kosch et al.*, 2002; *Pedersen et al.*, 2003]. The excitation energies $\varepsilon_\lambda$ of the $O(^1D)$ and $O(^1S)$ states responsible for the red and green lines are $\varepsilon_r$ = 1.96 eV and $\varepsilon_g$ = 4.17 eV, respectively. The population of energetic, $\varepsilon > \varepsilon_\lambda$, electrons can increase significantly due to stochastic and resonant interactions with plasma turbulence generated by a heating wave [e.g., *Gurevich et al.*, 1985; *Dimant et al.*, 1992; *Mantas and Carlson*, 1996; *Gurevich and Milikh*, 1997; *Istomin and Leyser*, 2003]. The former raises the electron temperature $T_e$, the latter accelerates a group of electrons in the high-energy tail of the initial distribution. Both effects are well documented [e.g., *Carlson et al.*, 1982; *Gustavsson et al.*, 2001].

[3] *Gustavsson et al.* [2002] emphasized that to interpret $\zeta_{gr} > 0.1$ in terms of electron heating requires unrealizable $T_e > 2$ eV, pointing out the importance of electron acceleration. It is commonly believed that electrons are most efficiently accelerated by Langmuir ($l$) turbulence. The generation of Langmuir waves is usually described in terms of nonlinear instabilities of ordinary ($o$) mode pump waves, known as the parametric decay (PDI) or oscillating two-stream (OTSI) instabilities [*Fejer*, 1979]. In the presence of background suprathermal ($\varepsilon \gg T_e$) electrons, many more energetic electrons are accelerated than would be in Maxwellian plasmas [*Mishin and Telegin*, 1986].

[4] Recent observations show that the airglow maximizes during HF injections toward magnetic zenith [*Kosch et al.*, 2002; *Pedersen et al.*, 2003]. The same is true for the intensity of Langmuir waves [*Isham et al.*, 1999] and electron heating [*Rietveld et al.*, 2003] observed by the EISCAT UHF radar. Furthermore, the red line is excited at magnetic zenith even at extremely low effective radiative power $P_0 \sim 2$ MW [*Pedersen et al.*, 2003].

[5] Ordinary mode waves with incident angles $\theta$ outside the Spitze region, $\theta > \theta_c$, reflect at altitudes $H_\theta$ below the standard reflection altitude $H_0$ where the local plasma frequency $f_p \simeq 10^4\sqrt{n_e}$ Hz equals the driver frequency $f_0$. Here $\theta_c = \arcsin(\sqrt{\frac{f_c}{f_c+f_0}}\sin\chi)$, $f_c$ is the local electron gyrofrequency, $n_e$ is the electron density in cm$^{-3}$, and $\chi$ is the magnetic dip angle [e.g., *Mjølhus*, 1990]. At HAARP $\chi \simeq 14.5°$ and $f_c \simeq 1.4$ MHz at altitudes near 200 km, so that for $f_0$ = 7 MHz the Spitze angle is $\theta_c \simeq 5.9°$.

[6] Figure 1 shows a schematic of ray trajectories for ordinary HF waves injected vertically and toward magnetic zenith into a horizontally stratified ionosphere at HAARP. Obliquely incident radiation does not form standing-wave patterns and swelling is absent. Thus, given $P_0$ = 150 MW and distance $R$ = 250 km, the wave amplitude is $E_0 \simeq 4.7\sqrt{P_0}/R \simeq 0.2$ V/m. With $f_0$ = 7 MHz and $T_e$ = 0.1 eV, the HF energy density at the reflection point is $W_0 = \frac{E_0^2}{4\pi} \simeq 10^{-4} n_e T_e$, sufficient to drive the PDI/OTSI at $H_0$ [*Fejer*,

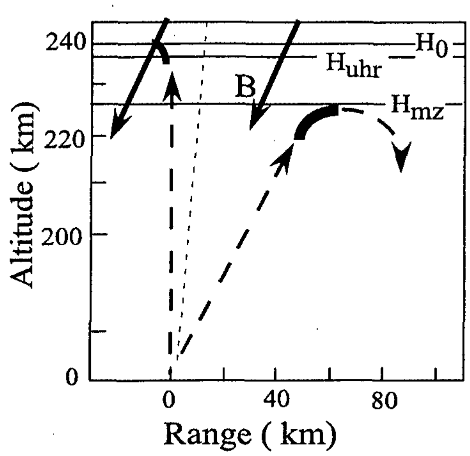


**Figure 1.** Schematic of ray propagation for ordinary HF waves injected toward local vertical and magnetic zenith. The magnetic field **B** direction is indicated by arrows. The heights of reflection and upper hybrid resonance are shown by horizontal lines. A light dashed line shows the Spitze angle direction. The regions of excitation of Langmuir turbulence appear as bold lines near the reflection points.

1979]. However, mismatch of frequencies at $H_\theta < H_0$ suppress these instabilities for injections outside the Spitze region.

[7] *Gurevich et al.* [2002] suggested that decreased plasma densities within striations generated by heating waves permit necessary phase matching [cf. *Muldrew*, 1978]. Changes in the refraction index (self-focusing) due to striations develop within tens of seconds after turn-on and explain some features of the spatial distribution of the HF-induced airglow. However, the rise time of Langmuir waves is ~10 ms [*Isham et al.*, 1999]. Besides, striations are generated in the upper hybrid layer [e.g., *Vaskov et al.*, 1981; *Lee and Kuo*, 1983]. Hence the reflection height of the heating wave must be at or above this height. For magnetic-zenith injections at HAARP, this can be satisfied only if $f_0 \leq 5.4$ MHz. However, the strongest airglow at HAARP occurred with higher $f_0$ values [*Pedersen et al.*, 2003].

[8] *Kuo et al.* [1997] showed that Langmuir waves can be excited by upper hybrid (*uh*) waves with amplitudes exceeding the threshold value $E^l_{uh} \sim 0.15$ V/m. Such *uh*-waves can be generated through the linear conversion of the *o*-mode on pre-existing field-aligned irregularities [*Wong et al.*, 1981] or through the parametric decay $o \rightarrow uh + lh$ [e.g., *Istomin and Leyser*, 1995]; *lh* stands for lower hybrid waves. Linear conversion within the upper hybrid layer proceeds with no set threshold. However, the parametric decay develops if $E_0 > E_0^{uh} \simeq 1.6 \tilde{f}_0^2$ mV/m, provided that the inequality $0.015 \lesssim |f_0 - s \cdot f_c| < 0.5$ MHz (integer $s \geq 3$) is valid. Here $\tilde{f}_0$ stands for the heating frequency in MHz. The rise time of the *uh*-wave is $\tau_{uh} \sim 1$–$3$ ms. Consistent with observations [*Isham et al.*, 1999], this makes the generation of Langmuir turbulence possible within ~10 ms.

[9] Besides electron impact, dissociative recombination (DR) of $O_2^+$ is an efficient source of the $O(^1D)$ and $O(^1S)$ states. Since the rate of dissociative recombination decreases with $T_e$, it is usually not considered in the theory of the HF-induced airglow. However, the quantum yield for the $O(^1S)$ state from dissociative recombination grows significantly with $T_e$ and with the vibrational temperature of oxygen ions [*Guberman*, 1997; *Peverall et al.*, 2001]. Furthermore, charge exchange $O^+ + O_2 \rightarrow O_2^+ + O$ is the dominant source of $O_2^+$ production in the nighttime *F* region. Its rate increases significantly whenever $O_2$ is vibrationally excited [*Viggiano and Williams*, 2001]. We expect a high degree of excitation of molecular species in the HF-illuminated region and thus enhanced yields of the green line emissions.

[10] A consistent theory of the HF-induced airglow accounting for the ion/neutral chemistry, modification of the heated spot, and self-focusing has yet to be worked out. We emphasize that responsible processes develop within tens of seconds. However, considerable airglow appears within first few seconds after the HF transmitter turns on. We designate this interval as the onset of HF-induced airglow.

[11] This paper develops a model for the onset of airglow enhancements at magnetic zenith accounting for the roles of $O_2^+$ dissociative recombination and energetic photoelectrons. The following section describes the onset characteristics from the HF heating experiments at HAARP. Section 3 describes the model. In particular, we propose a scenario (section 3.3) for the generation of strong Langmuir turbulence outside the Spitze region including magnetic zenith with upper hybrid waves as the primary source. The final section compares modeling results with optical measurements.

## 2. Airglow Onset: HAARP, February 2002

[12] The HAARP facility is located outside of Gakona Alaska (62.4°N, 145.15°W). During the course of HF heating experiments between 03:00 and 05:00 UT in February 2002 several passes of Defense Meteorological Satellite Program (DMSP) satellites flew to the east and west of the HAARP location. Each of these spacecraft carries a pair of upward looking particle spectrometers designed to measure fluxes of downcoming electrons and ions with energies between 30 eV and 30 keV. Consistent with prevailing quiet geomagnetic ($Kp = 2$) conditions, the equatorward boundary of auroral precipitation was several degrees in latitude poleward of Gakona. However, both before and after the local time of sunset, the spectrometers detected fluxes of Rdowncoming electrons whose spectra monotonically decreased with energies between 30 and 100 eV. These subauroral fluxes consist of photoelectrons that originated in the still sunlit southern ionosphere.

[13] Intense green-line emissions were observed between 03:48 and 05:00 UT (17:48–19:00 LT) on February 13, 2002 events of the HAARP optics campaign. Readers are referred to the report of *Pedersen et al.* [2003] for graphic examples of HAARP- induced airglow and the intensities of red/green-line emissions recorded during the period of interest. *O*-mode waves were injected toward magnetic zenith at $f_0 = 7.8$ MHz and at full power of 0.94 MW ($P_0 \simeq 165$ MW). The transmitter was programmed to turn on for 5 minutes exactly on the minute. Subsequent pulses followed 5 minutes pauses. However, the fourth pulse in the sequence started at 04:21:18 UT. A transmitter problem produced a false start, 04:30:00–04:31:00 UT, prior to the long pulse that began at 04:32:12 UT. In the course of this long pulse, the critical frequency of the *F* layer dropped

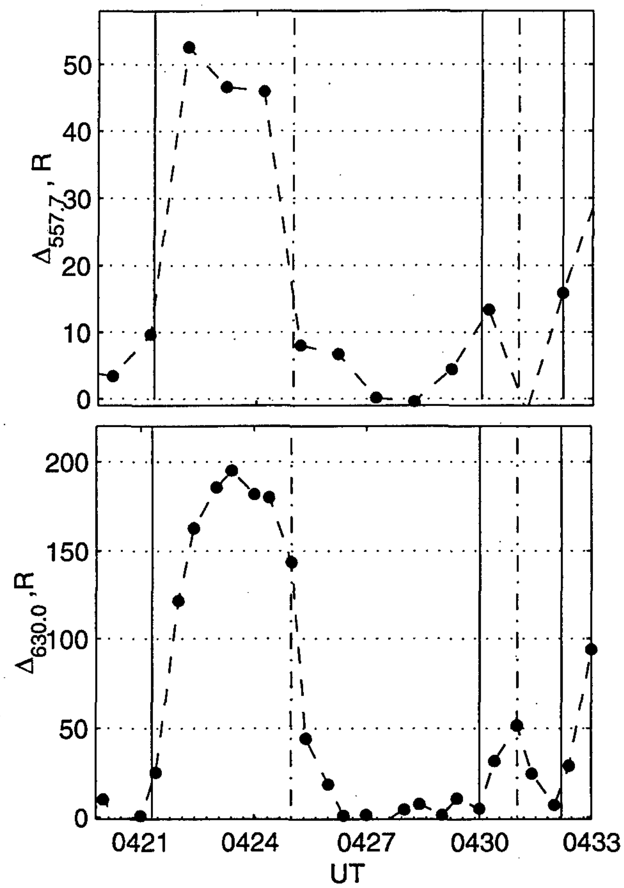


**Figure 2.** Heater-induced green (top panel) and red (bottom panel) lines during February 13, 2002 injections. Vertical solid and dash-dotted lines indicate heater turn-ons and turn-offs, respectively.

below the transmitter frequency and the sequence was terminated.

[14] To account for natural (background) variations, a polynomial fit was made to intensities measured along the magnetic meridian in each all-sky image, with the heated area blocked out. Throughout the studied period the difference between the background airglow $b_\lambda$ and that of the heated volume $I_\lambda$, during heater-off periods, was less than ~3–5 and ~7–10 R for green- and red-line emissions, respectively. This small difference determines the accuracy of heater-induced airglow measurements $\Delta_\lambda = I_\lambda - b_\lambda$.

[15] Figure 2 presents a detailed view of the last three HF-induced enhancements of February 13, 2002 at $f_0$ = 7.8 MHz. Data points are from two red line exposures compiled at 0 and 24 s after the beginning of each minute, and green-line exposure beginning at 12 s into the minute. After 04:00 UT, the exposure duration was 7.5 s. By examining cases when the heater turned on part way through an exposure, we see that the green-line emissions increased by ~5 (~10) R within ~1.5 (~7.5) s, while the red line gained ~25 (~30) R within ~13.5 (~20.5) s.

[16] Starting at 04:27:15 UT on February 9, 2002, a $P_0 \simeq$ 124 MW wave at a frequency 6.8 MHz was injected toward magnetic zenith. Figure 3 shows red- and green-line emissions before and during the pulse turn on from the heated volume and the background region. These data show that the difference between natural green-line emissions from magnetic zenith and ~70 km apart was less than 1 R. The same is true for the red-line emission before the heater turned on. This suggests that the background airglow was nearly uniform. When the heater turned on, green (red)-line intensities increased by ~5 (~30) R within ~4.5 (~16.5) s. It is relevant to note that *Kosch et al.* [2002] reported ~500 R enhancements of the red- and green-line emissions produced by the EISCAT superheater ($P_0$ ~ 600 MW) within the first 5-s frame.

## 3. Modeling Airglow Onset

[17] For modeling purposes we assume that the major species constituents of the neutral atmosphere and ionospheric plasma are well described by the MSIS-E [*Hedin*, 1991] and IRI [*Bilitza*, 2001] models, respectively. Although photodissociation of $O_2$ accounts for ~50% of the background airglow, it is unaffected by heating and hence is not included. Basic processes and their rate coefficients are discussed at length by *Solomon et al.* [1988], *Rees* [1989], *Witasse et al.* [1999], and *Fox and Sung* [2001]. Unless

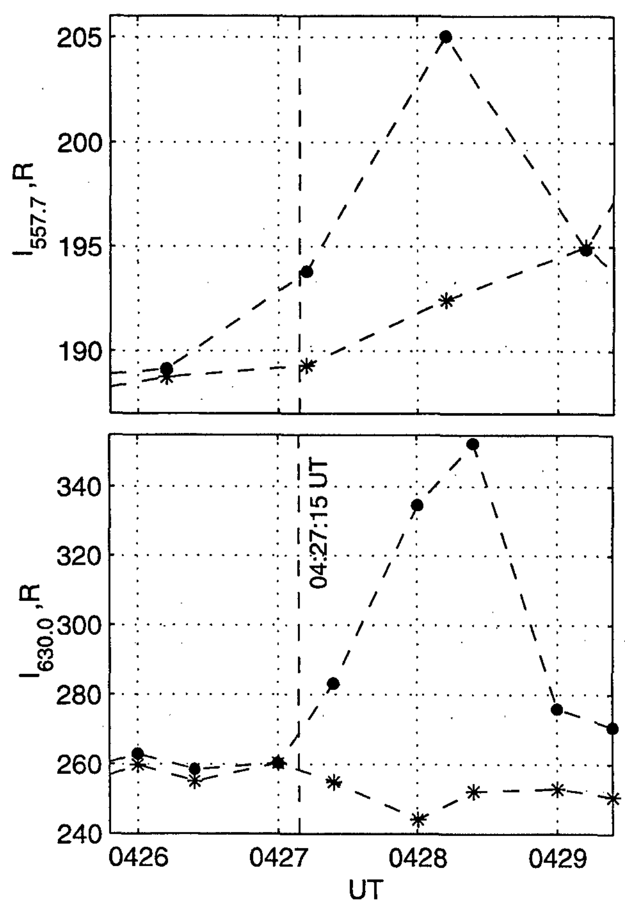


**Figure 3.** Heater-induced green (top) and red (bottom) lines during February 9, 2002 injections. Dots and asterisks stand for emissions from the magnetic zenith and background regions, respectively. Vertical dashed line indicates the heater turn-on time.

otherwise noted, we use the rate coefficients of *Fox and Sung* [2001].

[18] As mentioned above, radiative emissions from a heated volume reflect the superposed results of complex wave, plasma, and chemical interactions. To render our modeling of this complex chain of interactions intelligible, we divide this section into three parts. The first subsection concerns the basic contributors to the natural and artificially excited radiation budgets. The subsection on the background ionosphere estimates the distributions of ion species and energetic photoelectrons present at heater altitudes prior to turn on. The third subsection describes a flow chart encapsulating our concept of how heater-injected energy is transformed into the various wave modes that heat and/or accelerate ambient electrons that interact with ambient neutrals to produce onset green- and red-line emissions.

### 3.1. Basic Processes

[19] The red ($r$) and green ($g$) line-photons are emitted by atomic oxygen in the transition from the $O(^1D)$ and $O(^1S)$ states to $O(^3P)$ and $O(^1D)$ states, respectively. In photochemical equilibrium, the volume emission rate $\eta_\lambda$ is calculated from

$$\eta_\lambda^{eq} = A_\lambda \cdot [O_\lambda] = A_\lambda \frac{Q_\lambda}{L_\lambda + A_{\Sigma\lambda}} \quad (1)$$

Here $[O_\lambda]$ stands for the density of $O(^1D)$ or $O(^1S)$ in cm$^{-3}$; $Q_\lambda$ and $L_\lambda$ are the corresponding production and loss rates, respectively; $A_r \simeq 0.007$, $A_{\Sigma r} \simeq 0.009$, $A_g \simeq 1.2$, and $A_{\Sigma g} \simeq 1.3$ are the Einstein transition probabilities in s$^{-1}$. The column emission rate (in R) is found by integrating (1) along the line of sight

$$4\pi I_\lambda = 10^{-6} \int_{z_0}^{z_1} \eta_\lambda(z)\, dz \quad (2)$$

[20] The rates of electron impact excitation $O + e \rightarrow O(^1D, ^1S) + e$ and the Schumann-Runge dissociation $O_2 + e \rightarrow O + O(^1D) + e$ are calculated as follows

$$Q_\lambda^e = [O] \cdot \varkappa_\lambda^e = [O] \cdot 4\pi \int_{\varepsilon_\lambda}^{\infty} \sigma_\lambda(\varepsilon)\Phi(\varepsilon) d\varepsilon \quad (3)$$

Here $\Phi(\varepsilon)$ is the electron differential number flux, $\varkappa_\lambda^e$, $\varepsilon_\lambda$, and $\sigma_\lambda$ are the rate coefficients, threshold energies, and cross-sections, respectively. We employ electron impact cross-sections suggested by *Majeed and Strickland* [1997]. Figure 4 shows three of the major electron impact cross sections. It is worth noting that $\varkappa_r^e(T_e)$ and $\varkappa_g^e(T_e)$ calculated with a Maxwellian distribution of thermal electrons are close to the approximations used by *Mantas and Carlson* [1996] and *Gurevich and Milikh* [1997], respectively.

[21] The dissociative recombination of $O_2^+$ produces $Q_\lambda^{DR} = \varkappa_\lambda^{DR} n_e [O_2^+]$. The rate coefficients are usually approximated as $\varkappa_\lambda^{DR} = \beta_\lambda^{DR} \cdot \alpha$, where $\beta_\lambda^{DR}$ are the quantum yields. Estimated yields for $O(^1S)$ were subject to substantial discrepancy in the literature for many years. Recent experiments [*Kella et al.*, 1997; *Peverall et al.*, 2001; A. Petrignani et al., Vibrationally resolved rate coefficients and branching in the dissociative recombination of $O_2^+$, submitted to *Journal of Chemical Physics*, 2003] and theory

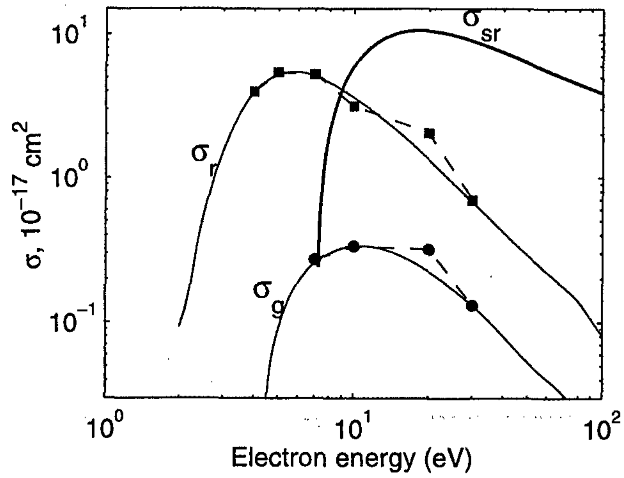


**Figure 4.** Cross-sections of the $O(^1D)$ and $O(^1S)$ excitation and the Schumann-Runge dissociation by direct electron impact. Squares and dots are the measured cross-sections from *Doering and Gulcicek* [1989]; solid lines are interpolations by *Majeed and Strickland* [1997].

[*Guberman*, 1988, 1997] appear to have reconciled this controversy. They established that $\alpha$ and $\beta_\lambda^{DR}$ depend not only on $T_e$ but also on the vibrational population $n_{O_2^+}^{(j)} = [O_2^+(j)]/[O_2^+]$ of $O_2^+$, where $j$ = 0, 1,... designates the vibrational level.

[22] For the ground vibrational state, $\beta_r^{(0)} \simeq 1.15$, $\alpha^{(0)} \simeq 2\ 10^{-7} \tau_e^{-0.63}$ [cf. *Fox and Sung*, 2001], and

$$\beta_g^{(0)} \simeq 0.063 \exp\left(-\frac{2.23}{\tau_e}\left(1 - \frac{0.39}{\tau_e}\right)\right) \quad (4)$$

for $1 \leq \tau_e \leq 10$, where $\tau_e = T_e$[K]/300. Importantly, the inequalities $\beta_g^{(1,2)} > \beta_g^{(0)}$ and $\beta_r^{(1,2)} < \beta_r^{(0)}$ hold. For the Maxwell-Boltzmann distribution of $N_{O_2^+}^{(j)}$ with vibrational temperatures $T_V \sim 0.25$ eV, one obtains $\beta_g^{DR} \sim 1.2\beta_g^{(0)}$ and $\beta_r^{DR} \sim 0.85\beta_r^{(0)}$.

[23] Collisional quenching $O^\# + X \rightarrow O + X$ is important for the $O(^1D)$ state, while for $O(^1S)$ at altitudes >150 km it can be ignored. Here $X$ stands for $O$, $O_2$, $N_2$, or $e_{th}$. The corresponding loss rates are $L_r^X = 10^{-11} \varkappa_X^q [X]$, where $\varkappa_O^q \simeq 0.65\ \tau_n^{0.14}$, $\varkappa_{O_2}^q \simeq 3.2 \exp(0.225/\tau_n)$, $\varkappa_{N_2}^q \simeq 1.8 \exp(0.36/\tau_n)$, and $\varkappa_e^q \simeq 28.7\ \tau_e^{0.91}$. Figure 5 shows the altitude profile of the total loss rate calculated with the MSIS-E and IRI parameters at 04:30 UT, February 13, 2002.

[24] Finally, the radiative cascade (RC) $O(^1S) \rightarrow O(^1D) + h\nu_{5577}$ yields $Q_r^{RC} = \eta_g$.

### 3.2. Background Ionosphere

[25] Simultaneous observations from a digisonde located at the HAARP site were used to correct the IRI model and to determine that the reflection height at magnetic zenith decreased from ~250 to ~235 km between 04:00 and 04:30 UT on February 13, 2002. At the same time, the shadow height increased from ~150 to ~230 km indicating the presence of photoelectrons [e.g., *Doering et al.*, 1975]. Furthermore, the southern-hemisphere region magnetically conjugate to HAARP remained in sunlight, even after local sunset. Thus a large fraction of the photoelectron spectrum

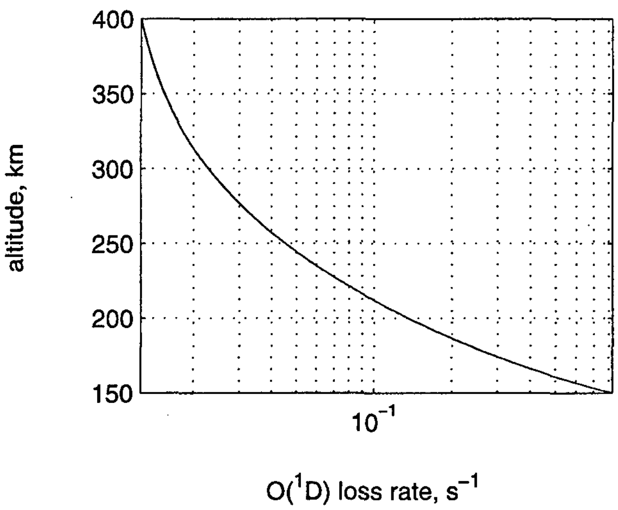


**Figure 5.** $O(^1D)$ loss rate.

generated in the conjugate ionosphere [e.g., *Peterson et al.*, 1977] had access to the ionosphere above HAARP. Electron spectrometers on DMSP satellites observed the high-energy tail of the conjugate photoelectron spectrum.

[26] Figure 6 shows altitude profiles of the background electron temperature $T_0$ and electron and $O^+$ densities calculated from the corrected IRI model. The twilight $O_2^+$ density above ~200 km is defined by the charge exchange and dissociative recombination

$$[O_2^+] \simeq \varkappa_{ce} \frac{[O_2][O^+]}{\alpha n_e} \tag{5}$$

Here the rate of charge exchange $\varkappa_{ce}$ after *Viggiano and Williams* [2001] is

$$\varkappa_{ce} \simeq \varkappa_{ce}^{(0)}\left(n_{O_2}^{(0)} + n_{O_2}^{(1)}[4.9\lg(\tau_{in}+1) - 2.5]\right)$$
$$\varkappa_{ce}^{(0)} \simeq 10^{-11}\left(\frac{2.2}{\tau_{in}^{0.52}} + 6.05\exp\left(-\frac{34.04}{\tau_{in}}\right)\right), \tag{6}$$

$\alpha$ is the recombination rate in cm$^3$/s, $n_{O_2}^{(j)} = [O_2(j)]/[O_2]$, $[Y]$ stands for the density of species $Y$ in cm$^{-3}$, and $\tau_{in} = 1.5\tau_i + \tau_n - 1$ with $\tau_{i,n} = T_{i,n}$[K]/300.

[27] The IRI model predicts $T_e = T_0 \leq 0.2$ eV (Figure 6). At these temperatures and $n_e > 10^5$ cm$^{-3}$, a Maxwellian distribution $F_M$ is a good representation of the thermal electron distribution function ($F_{th}$) in the $F$ region [*Mishin et al.*, 2000]. The differential number flux $\Phi_s(\varepsilon)$ of photoelectrons in the $F$ region [e.g., *Rees*, 1989] at $\varepsilon \geq \varepsilon_s = \frac{1}{2}mv_s^2 \sim 5$ eV can be approximated by a power-law function

$$\Phi_s(\varepsilon) = \frac{p_s - 0.5}{4\pi\varepsilon_s} n_s v_s \left(\frac{\varepsilon_s}{\varepsilon}\right)^{p_s} \tag{7}$$

where $n_s \leq 100$ cm$^{-3}$ and $p_s \simeq 3$.

### 3.3. HF-Perturbed Ionosphere

[28] Our scenario for the excitation of plasma turbulence and subsequent electron energization in the HF-illuminated region at magnetic zenith is represented schematically in Figure 7. First, upper hybrid waves are generated due to the parametric decay $o \rightarrow uh + lh$. Its threshold field is $E_0^{uh} \sim 0.08$ V/m if the matching conditions are met at $x_{uh} = k_{uh}^2 r_e^2 < 1$, where $r_e$ is the thermal electron gyroradius and $k_{uh}$ is the $uh$-wave vector. Landau damping of short-scale lower hybrid waves raises the threshold otherwise [*Mishin et al.*, 1997]. For $f_0$ = 7.8 (6.8) MHz and $\theta \simeq 14.5°$, the matching conditions at the reflection height $H_0$ are easily satisfied for $x_{uh} \simeq 0.3$ ($\simeq$0.15). Finally, the free space field of the incident wave $E_0 \simeq 0.25$ (0.2) V/m exceeds $E_0^{uh}$.

[29] When the amplitude of the primary $uh$-wave $E_{uh}$ exceeds ~10 mV/m, it excites other, lower frequency, $uh$-wave with the growth rate of order $\gamma_{uh} \sim \tau_{uh}^{-1}$ [*Zhou et al.*, 1994]. The same is true for subsequently generated $uh$-waves. The frequency-step of this spectral transfer is quite small $|\delta\omega_{uh}|/\omega_{uh} \sim 0.1 x_{uh}\delta k_{uh}/k_{uh} \sim 0.2\sqrt{m_e/m_i}$, where $m_e/m_i$ represents the electron/ion mass ratio. As many spectral steps occur before the parametric instability saturates, the resulting $uh$-energy spectrum consists of a large number $\Lambda \gtrsim 10$ of spectral peaks, each of order $E_0^2/(\gamma_{uh}\tau_{uh})$ (well-known weak turbulence "cascading" [e.g., *Sagdeev et al.*, 1991]). Thus the total $uh$-energy density can be estimated as $W_*/\Lambda \sim W_0 \simeq 7.5$ (5) eV/cm$^3$ with the r.m.s amplitude $E_* \simeq \sqrt{4\pi W_*} \sim \sqrt{\Lambda}E_0$.

[30] Upper hybrid waves with amplitudes $E_{uh}$ exceeding the threshold value $E_{uh}^l \simeq 0.1/x_{uh}^{1/2}$ V/m generate short-scale, $k_l > 3^{-1/2}/r_e$, Langmuir waves [*Kuo et al.*, 1997]. Due to a weak dependence of the growth rate $\gamma_l$ on $k_{uh}$, one can take $E_{uh} \sim E_*$. As this value greatly exceeds the threshold $E_{uh}^l$, the growth rate can be evaluated as $\gamma_l \sim \nu_e E_{uh}/E_{uh}^l$ [*Kuo et al.*, 1997], which amounts to $\gamma_l \sim \nu_e\sqrt{10^2 x_{uh}\Lambda E_0^2}$. Here $\nu_e$ is the electron elastic (transport) collision frequency.

[31] The growth of Langmuir waves at $T_e/T_i \leq 4$ saturates via induced scattering by ions that transfers energy toward

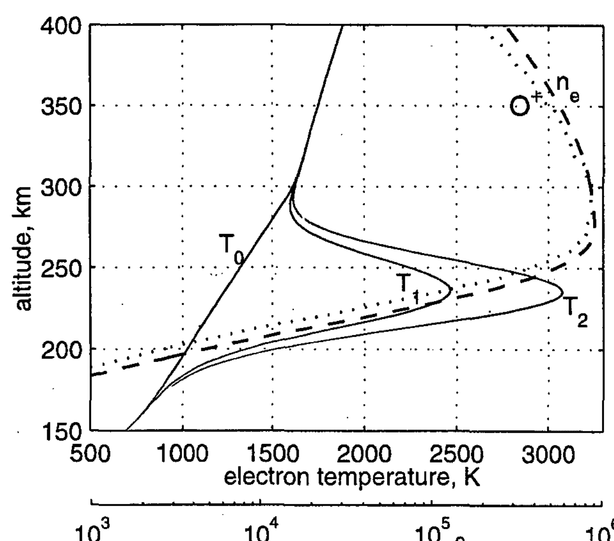


**Figure 6.** Altitude profiles of the electron temperature (solid lines), electron density (dashed line), and $O^+$ density (dotted line). $T_0$, $T_1$, and $T_2$ stand for 0, 1, and $\simeq$1.5 s after turn-on.

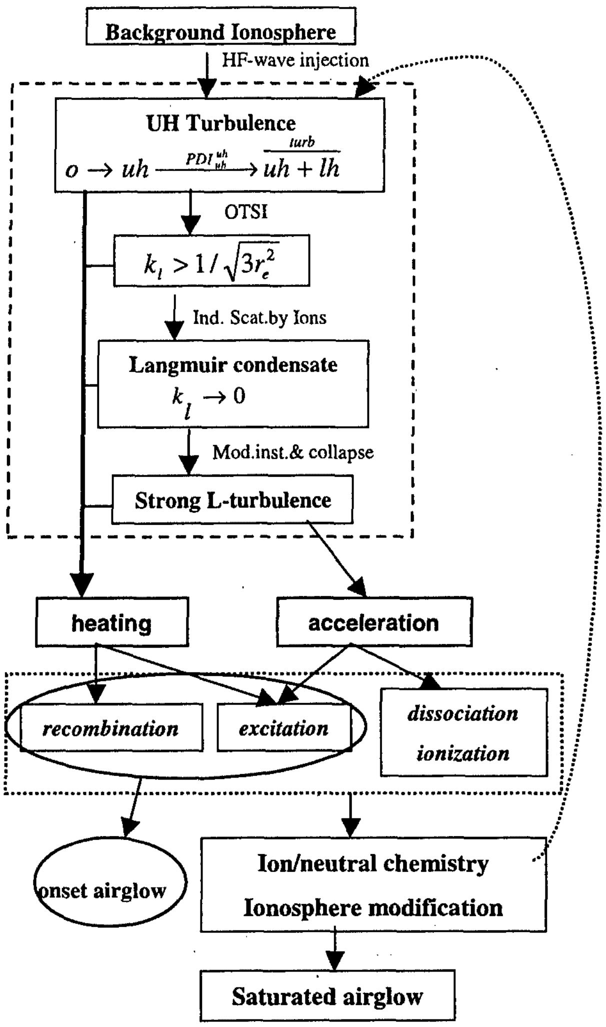


**Figure 7.** Block-diagram showing synergy of a model of HF-induced airglow at magnetic zenith. Shaded boxes indicate the steps to be worked out.

small $k$ [e.g., *Zakharov et al.*, 1976]. This process is governed by the pump $uh$-wave power. In particular, if the width $\Delta k_l$ of the parametrically unstable $l$-wave spectrum exceeds $\Delta k_c \simeq k_l\nu_e/\gamma_l$, the Langmuir wave energy $W_l$ accrues in a state with $k \rightarrow 0$ (Langmuir condensate) [*Zakharov et al.*, 1976; *Zakharov*, 1984]. The spectral width can be evaluated as $\Delta k_l \sim 3(f_0/f_c)^2 W_*/(n_eT_er_e)$. Thus, for $E_0 > 0.15(\tilde{f}_0\nu_e)^{1/6}/\sqrt{\Lambda} \lesssim 0.2$ V/m the "condensate" condition $\Delta k_l > k_l\nu_e/\gamma_l$ is fulfilled. The dynamics of the Langmuir condensate is defined by the modulational instability and collapse leading to establishment of strong (cavitating) Langmuir turbulence [e.g., *Galeev et al.*, 1977; *Zakharov*, 1972, 1984].

[32] We emphasize that this applies not only at the reflection layer $H_0$ but also well below it, wherever the matching conditions for the parametric decay $o \rightarrow uh + lh$ are met and local values of $x_{uh} < 1$. Given the plasma density profile shown in Figure 6 and $f_0 \sim 6$–8 MHz, the altitude extent of this region is ~10 km. One should compare this value with a few 100-m size of the Airy pattern at $H_0$, where the PDI/OTSI instabilities develop during injections within the Spitze region [*Fejer*, 1979]. We believe that this is the key factor in exciting strong airglow at magnetic zenith.

### 3.3.1. Heating

[33] Collisional damping of high-frequency $o$-, $uh$-, and $l$-waves is the major source of (stochastic) electron heating. *Mishin et al.* [2000] showed that at $[N_2] > 10^7\tilde{f}_0^2$ cm$^{-3}$ the process of $N_2$-vibrational excitation dominates the formation of the distribution function of ionospheric electrons with energies $3.5 \geq \varepsilon > \varepsilon_{vib} \simeq 1.8$ eV, which can be represented as follows

$$F_{th}(\varepsilon > \varepsilon_{vib}) \simeq F_M\left(\frac{\varepsilon + \varepsilon_{vib}}{2}\right)\sqrt{\frac{\varepsilon^{0.5}\kappa(\varepsilon_{vib})}{\varepsilon_{vib}^{0.5}\kappa(\varepsilon)}} \times \exp\left[-\frac{1}{\sqrt{2T_e}}\int_{\varepsilon_{vib}}^{\varepsilon}\xi^{-1/2}\kappa(\xi)d\xi\right] \quad (8)$$

Here $\kappa^2(\varepsilon) = \nu_{il}(\varepsilon)/\nu_{ee}(\varepsilon) + 0.5\varepsilon/T_e$, where $\nu_{il}/\nu_{ee}$ is the ratio between electron inelastic and electron-electron collision frequencies. A total HF wave energy density $W_{HF} \ll 10^{-2}n_0T_e$ is assumed. At energies $\varepsilon < \varepsilon_{vib}$ the distribution is close to $F_M$, while at $\varepsilon > 3.5$ eV it may slightly deviate from $F_M$ due to the $O(^1D)$ and $O(^1S)$ excitation.

[34] The electron temperature can be evaluated from the energy balance

$$1.5n_e\partial T_e/\partial t = \Gamma_e - \nu_e\delta \cdot n_eT_e - \nabla_\| q_{e\|} \quad (9)$$

where $q_{e\|}$ is the (parallel) electron heat flux, $\Gamma_e \simeq \nu_eW_{HF}$ is the volume heating rate, and $\delta = \langle\nu_{il}/\nu_e\rangle$ is the coefficient of inelastic losses averaged over the total distribution $F_{th}$.

[35] Given $\nu_e \sim 500$ s$^{-1}$, $\Gamma_e$ is $\gtrsim 40$ (25) keV/cm$^3$/s. Note that to match their optical observations, *Mantas and Carlson* [1996] used the values of $\Gamma_e$ between $\simeq 50$ and $\simeq 125$ keV/cm$^3$/s at ~260 km for $f_0 \simeq 5$ MHz. *Gustavsson et al.* [2001] used $\Gamma_e \simeq 60$ keV/cm$^3$/s for $f_0 \simeq 4$ MHz to fit the electron temperature profile maximum 4000 K at ~220 km observed by the EISCAT UHF radar. Neither of them accounted for the decrease of $\delta(T_e)$ due to the deviation of the thermal electron distribution from Maxwellian. Taking that into account, the calculation results of *Mantas and Carlson* [1996] and *Gustavsson et al.* [2001] can be scaled to the $T_{1,2}$-profiles in Figure 6, pertaining to our case.

### 3.3.2. Acceleration

[36] Resonant $lh$- and $l$-wave-particle interaction accelerates electrons. *Musher et al.* [1978, 1986] analyzed in detail the dynamics of $lh$-waves excited by an external source. At $1 < T_e/T_i \leq 4$, high-frequency waves, $f_{lh} > f_c\sqrt{2m_e/m_i}$, are dominated by induced scattering by ions. The rate of spectral transfer toward smaller frequencies is $\tilde{\gamma}_{lh} \sim (f_c^2/f_{lh})W_{lh}/n_0T_e$, where $W_{lh}$ is the energy density of lower hybrid waves. Equating the growth rate $\gamma_{uh}$ to $\tilde{\gamma}_{lh}$ yields the energy density at the saturated state $W_{lh}^0/n_0T_e \sim \gamma_{uh}f_{lh}/f_c^2 \sim 10^{-5.5}$.

[37] The dynamics of $lh$-waves in the low-frequency, $f_{lh} < f_c\sqrt{2m_e/m_i}$, region is dominated by the lower hybrid collapse [e.g., *Musher et al.*, 1978]. The threshold energy density is quite low $W_{lh}^c/nT_e \simeq x_{lh}(m_e/m_i)(f_c/f_0)^2 \leq 10^{-6}$, and is surely exceeded in our case. In the course of collapse, the longitudinal and transverse dimensions of a $lh$-cavity,

$l_\parallel \sim l_\perp \sqrt{m_i/x_{lh} m_e}$, diminish. Finally, the *lh*-wave energy in a collapsing cavity is absorbed by particles via Landau damping [*Musher et al.*, 1978, 1986; *Sotnikov et al.*, 1978; *Shapiro et al.*, 1993]. The density of accelerated electrons in the magnetic field-aligned high-energy tail of the distribution function, determined by assuming resonant particles carry away all of the energy pumped into collapsing cavities [e.g., *Shapiro et al.*, 1993], is of order $<10^{-5} n_e$. This *lh*-contribution appears insufficient to account for the observed HF-induced airglow. It is worth noting that *lh*-collapse also leads to transversely accelerated ions that absorb the collapsing energy at approximately the same rate as electrons [*Sotnikov et al.*, 1978].

[38] Strong Langmuir turbulence consists of an ensemble of collapsing cavitons that transfer the wave energy toward small scales. Small-scale, $k \geqslant \omega_k/v_{\min}$, waves are absorbed by fast $v \geqslant v_{\min} \gg \sqrt{T_e/m_e}$ electrons via Landau damping, thereby increasing the population of suprathermal electrons. The distribution of accelerated electrons $F_a(\varepsilon) = \frac{m^2}{2\varepsilon}\Phi_a(\varepsilon)$ in weakly magnetized $f_c \ll f_p$ plasmas can be found from the kinetic equation [e.g., *Galeev et al.*, 1977]

$$\frac{\partial F_a}{\partial t} = \frac{1}{v^2}\frac{\partial}{\partial v}\left[\frac{\omega_k^4}{mnv}\int_{\omega_k/v}^{\infty}\frac{W_k}{k^3}dk\frac{\partial F_a}{\partial v}\right] \quad (10)$$

where $\omega_k/(2\pi) \sim f_p$ and $W_k$ are, respectively, the frequency and spectral energy density of Langmuir waves. The latter is determined from the requirement for energy balance

$$\frac{dW_k}{dt} + \frac{d}{dk}\left[W_k\frac{dk}{dt}\right] = \Gamma_k W_k \quad (11)$$

Here

$$\Gamma_k = -\left(\frac{2\pi^2}{n_0}\right)\left(\frac{\omega_k^4}{k^3}\right)F_a\left(\frac{\omega_k}{k}\right) \quad (12)$$

is the rate of Landau damping; $k(t) \sim (t_c - t)^{-2}$ is defined by the collapse law in the absorption region [*Pelleiter*, 1982], and $t_c$ is the collapse time.

[39] The steady state solution of (10)–(12) at $\varepsilon \geq \varepsilon_{\min} = \frac{1}{2}m_e v_{\min}^2$ is a power-law function [*Pelleiter*, 1982; *Shapiro and Shevchenko*, 1984] that yields the differential number flux

$$\Phi_a(\varepsilon) = \frac{p_a - 0.5}{4\pi\varepsilon_{\min}} n_a v_{\min}\left(\frac{\varepsilon_{\min}}{\varepsilon}\right)^{p_a} \quad (13)$$

($p_a \simeq 0.75$). One-dimensional (magnetically field-aligned) numerical modeling of electron acceleration by strong Langmuir turbulence yields $\Phi_a(\varepsilon_\parallel) \sim \varepsilon_\parallel^{-1}$, consistent with the 1D scaling law for Langmuir collapse [*Galeev et al.*, 1983; *Wang et al.*, 1997]. A flat distribution of accelerated electrons is consistent with the observations of *Carlson et al.* [1982].

[40] The minimum energy $\varepsilon_{\min}$ and the density of particles in the high-energy tail $n_a$ are determined by the boundary condition

$$\Phi_a(\varepsilon_{\min}) = \Phi_0(\varepsilon_{\min}) \quad (14)$$

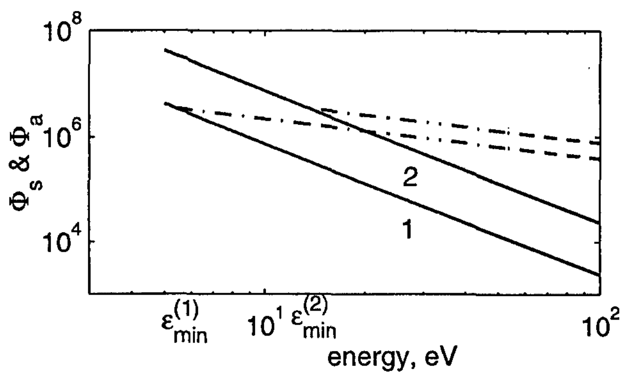


**Figure 8.** Differential number fluxes of photoelectrons (solid lines) and accelerated electrons (dash-dotted lines) in $cm^{-2}s^{-1}ster^{-1}eV^{-1}$ for $n_s = 1$ (line 1) and 10 (line 2) $cm^{-3}$.

and the wave energy flux transferred by collapsing cavitons and absorbed by accelerated particles

$$\omega_p\left(\frac{m_e}{m_i}\frac{|E_{k_m}^2|}{24\pi n_0 T_e}\right) \simeq -\Gamma_{k_m} \quad (15)$$

Here $E_{k_m}$ is the wave amplitude in a cavity of the scale length $l_{\min} \sim k_m^{-1} \simeq v_{\min}/\omega_p$ when collapse is arrested and $\Phi_0$ is the differential number flux of ionospheric electrons.

[41] For a Maxwellian electron distribution, one has $\varepsilon_{\min}^M \sim 20T_e$. When photoelectrons are present, $\Phi_0(\varepsilon) \to \Phi_s(\varepsilon)$ at $\varepsilon \gg T_e$, and $\varepsilon_{\min}$ is defined by

$$\varepsilon_{\min}^s \sim 30(n_s T_e/W_l)^{2/5}\ \mathrm{eV} \quad (16)$$

valid for the energy density of Langmuir turbulence $W_l \gg n_0 T_e m_e/m_i$ [*Mishin and Telegin*, 1986]. Given $n_s = 10\ cm^{-3}$, $n_e = 6\ 10^5\ cm^{-3}$, and $W_l \simeq 10^{-4} n_e T_e$, from (16) one gets $\varepsilon_{\min}^s \simeq 15$ eV. The number of accelerated electrons exceeds the photoelectron background in the energy range $\varepsilon > \varepsilon_{\min}^s$ by a factor of $\frac{p_s - 0.5}{p_a - 0.5} \sim 10$. For energies $\varepsilon \leq \varepsilon_{\min}^s$, the distribution is unaffected. Figure 8 shows differential number fluxes of photoelectrons (7) and accelerated electrons (13) for $n_s = 1$ and 10 $cm^{-3}$, $n_e = 6\ 10^5\ cm^{-3}$, and $W_l \simeq 10^{-4} n_e T_e$.

## 4. Discussion

[42] Figure 9 shows the components of the equilibrium volume emission rates $\eta_X^{eq}(1)$ for the period near 0430 UT. To calculate the contributions of photoelectrons (7) and accelerated electrons (13), $n_s = 10\ cm^{-3}$ and $W_l \simeq 10^{-4} n_e T_2$ were chosen, yielding $\varepsilon_{\min}^s \simeq 15$ eV for $f_0 = 7.8$ MHz. As expected from (4), the dissociative recombination contribution to the red line is reduced much more than that for the green line. Apparently, heated electrons dominate the red-line emission, while accelerated electrons are the major contributor to the green line. Note a factor of about 2 difference between calculations with $F_M(T_1)$ and $F_{th}(T_1)$ consistent with *Mishin et al.* [2000] (it is $\simeq 3.5$ for $T_e = T_2$).

[43] In the heated volume the densities of excited oxygen atoms and of $O_2^+$ ions grow. During the onset period $[O(^1D)]$ and $[O_2^+]$ remain well below the new equilibrium values that

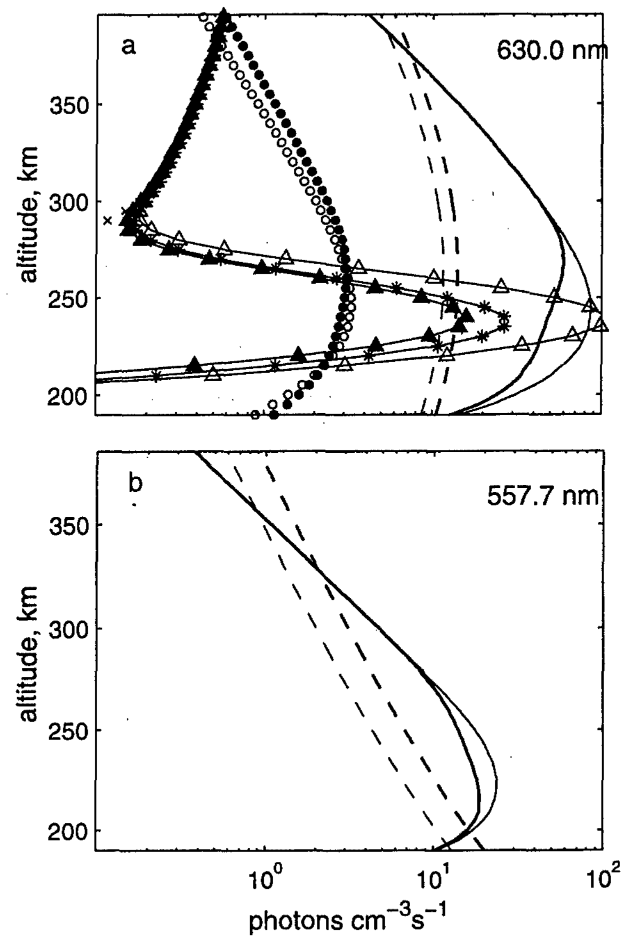


**Figure 9.** Equilibrium volume emission rate components of the red (a) and green (b) emissions: Dissociative recombination with $[O_2^+]_0$ from the background $T_e = T_0$ (thin solid line) and heated $T_e = T_2$ (bold solid line) regions; photoelectrons (thin dashed line) and accelerated electrons (bold dashed line); $O(^1S)$ cascading from the background (circles) and heated (dots) regions; thermal electrons with $T_e = T_0$ (crosses), $T_1$ (filled triangles), and $T_2$ (open triangles); Maxwellian electrons with $T_e = T_1$ (asterisks).

are of the order of $Q_r/L_r$ and $\varkappa_{ce}$ $[O_2]/\alpha(T_2)$, respectively. Corresponding time-variations of the contributing components to the total volume emission rate during the onset can be evaluated assuming that $A_{\Sigma g} < \zeta < \min[1/(n_e\alpha_2), 1/L_r] \simeq$ 15 s and $L_r \simeq \mathrm{const}(t)$ at 200–300 km. Here $\zeta = t - t_0$ and the heating time $t_0 \sim 1$ s. As a result, one gets

$$\delta[O_2^+] = [O_2^+]/[O_2^+]_0 - 1 \simeq n_e[\alpha(T_0) - \alpha(T_2)] \cdot \zeta \quad (17)$$

This yields

$$\eta_g(\zeta) \simeq \eta_g^{eq} + \eta_g^{DR} n_e[\alpha(T_0) - \alpha(T_2)] \cdot \zeta$$
$$\eta_r(\zeta) \simeq (\eta_r^{eq} L_r + 0.5\eta_r^{DR}[L_r + n_e\alpha(T_2)]\delta[O_2^+])\zeta \quad (18)$$

[44] Using (18) and Figure 9, the variation of the red- and green-line emissions during the onset period are calculated from (2). The results are shown in Figure 10 together with the observed intensities (section 2). There is good apparent agreement with the HAARP observations.

[45] We emphasize that the onset values of the HF-induced airglow amount to $\lesssim 10\%$ of the saturated values. As the reflection heights of 7.8 and 6.8 MHz waves at magnetic zenith are below the upper hybrid resonance, similar to Figure 1, striations are not generated. Thus self-focusing due to striations seems unlikely. This indicates that the ionosphere's parameters in the heated volume are significantly modified. Apparently, a complete explanation of artificially induced airglow involves a complex chain of plasma and chemical interactions within the heated volume and has yet to be worked out.

## 5. Conclusion

[46] Observations of airglow at 630 nm (red line) and 557.7 nm (green line) acquired during the February 2002 campaign at the HAARP heating facility have been analyzed. During injections toward magnetic zenith the intensities of green and red line emissions gain ~5–10 R within ~1 s and ~20 R within ~10 s, respectively. We call this period the onset of HF-induced airglow and develop an explanatory model for the growth of emissions from magnetic zenith. The model accounts for the roles played by ambient photoelectrons and dissociative recombination of $O^+$ ions and shows that heating and acceleration of background electrons, possibly by strong Langmuir turbulence, dominate the airglow onset.

[47] We also propose a new scenario for the generation of strong Langmuir turbulence for injections outside the Spitze region, including magnetic zenith, with upper hybrid waves as the primary source of Langmuir waves. Those waves accumulate in the region of small wavevectors due to induced scattering by ions (Langmuir condensate). Waves in the condensate are subject to the modulational instability and collapse thus leading to strong turbulence. This scenario

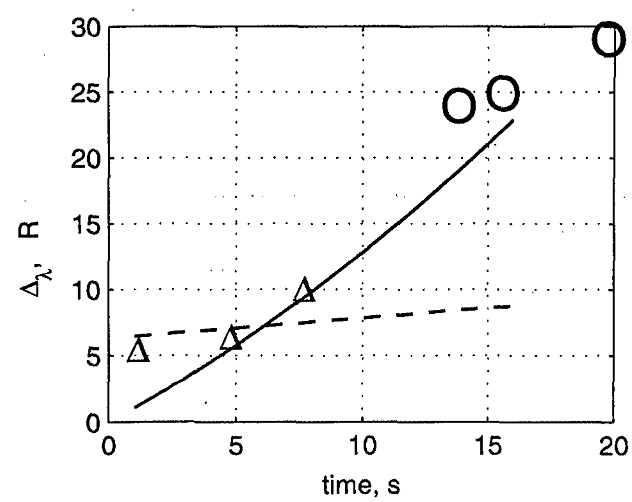


**Figure 10.** Calculated intensities of the HF-induced red (solid line) and green (dashed line) emissions during the onset period. The observed intensities are shown by open circles (red) and triangles (green), respectively.

applies not only at the reflection layer but also well below it, wherever the matching conditions for the parametric decay $o \rightarrow uh + lh$ are met and local values of $x_{uh} < 1$. For injections toward magnetic zenith and heating frequencies $\geq$5.4 MHz, this occurs below the altitude of the upper hybrid resonance. The large altitude extent of the strong turbulence region appears to be the key factor in exciting strong airglow at magnetic zenith.

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**14. ABSTRACT**

[1] Observations of airglow at 630 nm (red line) and 557.7 nm (green line) during the February 2002 campaign at the High Frequency Active Auroral Research Program (HAARP) heating facility are analyzed. We find that during injections toward magnetic zenith (MZ) the green and red lines gain ~5 R within ~1 s and ~20 R within ~10 s, respectively. We term this period the onset of the HF-induced airglow. A model of the onset at magnetic zenith is developed. It accounts for background photoelectrons and dissociative recombination of $O_2^+$. It is shown that heating and acceleration of background electrons dominate the airglow onset. We propose a scenario for the generation of strong Langmuir turbulence for injections outside the Spitze region, including magnetic zenith.

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