

Coherent plasma processes and the continuum emission in active galactic nuclei

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SUMMARY

Stimulated Raman scattering processes appear to be able to produce the typical electromagnetic continuum of an active galactic nucleus (AGN). Seed photons beat through Raman forward-scattering to create Langmuir plasma waves which can quickly accelerate electrons to Lorentz factors of 10^3 to 10^4 . These electrons radiate their energy through Raman back-scattering off magnetic fields engendered by magnetic modulational instabilities afflicting the Langmuir waves. The frequency of emission is proportional to the square root of the ambient density, so γ -rays can be produced in the highest density regions near the central engine, with X-rays through to IR photons produced at greater distances. Both the high luminosity and broken power-law continuum characteristic of an AGN emerge from reasonable density distributions. Bremsstrahlung emission in the UV is a by-product of this mechanism and could explain the ‘blue bump’. Because the electrons are continuously accelerated in this picture, a steady-state distribution arises naturally and a simple relation exists between the density of the relativistic particles which emit the radiation and the ambient plasma density which is involved in the acceleration.

1 INTRODUCTION

The significant body of evidence that most Seyfert galaxies and quasars have very similar continuum spectra from the infrared through X-ray bands has naturally attracted much theoretical attention. This so-called ‘universal’ spectrum for AGN has $\alpha_{\text{IX}} \approx 1.2$ from the IR to far UV, and $\alpha_{\text{X}} \approx 0.7$ in the X-ray (with $f_{\nu} \propto \nu^{-\alpha}$), and now has considerable observational backing (Lawrence & Elvis 1982; Tananbaum *et al.* 1983; Rothschild *et al.* 1983; Malkan 1984; Petre *et al.* 1984; Worrall & Marshall 1984; Stein & O’Dell 1985; Ward *et al.* 1987; Mushotzky 1988). Still, it must be noted that there is significant variation in the X-ray spectra of quasars (Elvis & Lawrence 1985; Wilkes & Elvis 1987) and a sum of log-normal (Gaussian) components might provide a superior fit, at least in some cases (Perry, Ward & Jones 1987). Although few observations have penetrated the extreme UV band, there is typically some break in the spectrum where these two continua meet between 12 and 750 eV (Petre *et al.* 1984; Elvis, Wilkes & Tananbaum 1985; Mushotzky 1988). It is possible that a third component contributes in this area, producing the measured variations in the X-ray spectra with

an underlying $\alpha_{\text{X}} \approx 0.5$ (Wilkes & Elvis 1987). Excess emission in the IR is also seen in many AGN, and much of the IR could be dust-reprocessed UV emission (e.g. Carleton *et al.* 1987), but this conclusion is far from established, especially for Seyfert 1s and quasars (e.g. O’Brien, Gondhalekar & Wilson 1988).

A surprisingly large number of plausible explanations for the origin of this continuum have been put forward. In general, the favoured pictures have involved incoherent radiation mechanisms. The most common is synchrotron self-Compton emission of non-thermal electrons, which seems to work rather well from the IR to UV, particularly for blazars (e.g. Stein & O’Dell 1985, Stein 1988). But Comptonized, self-absorbed thermal cyclotron radiation from mildly relativistic electrons also appears to be able to fit that part of the spectrum (e.g. Begelman 1988). The nature of the higher energy part of the spectrum has been much more vigorously debated; many of the earlier proposals were reviewed by Mushotzky (1984). A real problem has been to describe how very high-energy electrons can be accelerated in AGN. One possibility is direct acceleration of electrons in accretion flows or shocks near a black hole (e.g. Cowsik &

Lee 1982; Ellison & Eichler 1984). More promising models involve the indirect production of relativistic electrons from relativistic protons. One indirect approach uses inelastic proton-proton collisions with pion decays yielding numerous secondary electrons with good fits to the overall spectrum (Kazanas & Ellison 1986; Zdziarski 1986; Dermer 1988). An interesting variation on this theme uses proton-photon interactions to yield an electron-positron plasma, either through direct pair creation or through photo-produced pions (Sikora *et al.* 1987; Begelman 1988).

In this paper we discuss a significantly different mechanism that should be added to the list of continuum production candidates: coherent plasma processes, in particular, a combination of stimulated Raman scattering processes. In this picture, the gist of which was first proposed by one of us (Krishan 1983, 1985), and which has been summarized in Krishan & Wiita (1986), a region of suitably varying density containing fast particles and electromagnetic fields can generate the observed continuum all the way from the IR (and even radio) through hard X-rays (and even γ -rays). We note that electrodynamic processes allow particles to be in a state of continuous excitation since they can be accelerated by the fields that they themselves emit (*cf.* Pacini & Salvati 1982).

The essence of our model is the following: electrons are accelerated by Langmuir waves which are generated by Raman forward-scattering (RFS) in the ambient plasma. Simultaneously, spatially periodic magnetic fields and/or low-frequency electromagnetic waves (LFEW) are produced by magnetic modulational instabilities of the Langmuir field. The bulk of the observed continuum radiation is emitted via Raman back-scattering (RBS) off the periodic magnetic fields or LFEW. Our scenario requires the concerted action of several plasma processes, but, as we shall demonstrate, it is very reasonable to expect that such processes will arise in the region surrounding the supermassive black hole (SMBH), assumed to provide the ultimate energy source for AGN. The emission region decreases in density with distance from the SMBH as does the emission frequency, so that the highest energy photons are produced at a few Schwarzschild radii (R_s), and the IR from significantly greater radii (but still less than a parsec). Related plasma processes have recently been shown to be important for anomalous heating of emission-line regions (Krishan 1987) even if the standard values of gas parameters based upon incoherent processes are employed.

Our scenario has some very nice properties that distinguish it from most other proposed generators of the continuum. First, the acceleration and radiation processes are simultaneous and complementary; they naturally produce steady-state values for the fast particles' energies, essentially regardless of the injection spectrum. This also means that the same group of particles can give rise to a wide range of photon energies as they move outwards into regions of lower density. Relatively minor changes in the density distributions can produce the observed break in the sub-keV region, and variability at different bands can arise from quasi-coherent plasma density fluctuations naturally occurring in magnetohydrodynamical flows.

An early attempt to apply plasma processes to quasar continua is that of Colgate, Lee & Rosenbluth (1970), whose picture is in some respect the opposite of ours. They argued

that strong plasma turbulence allows Langmuir-type waves to be engendered by streaming ejecta; those waves can yield photons at twice the plasma frequency through non-linear interactions. Those photons can then be scattered off plasma waves, executing a random walk in phase space and allowing some to reach energies adequate to explain the radio emission. The coherence of the radiation in our picture indicates that it ought to dominate, if the conditions are such that the competing processes can both function, and we note that our mechanism can yield the bulk of the entire continuum. Recently Baker *et al.* (1988) have considered an alternative approach to coherent processes and have shown they have great promise when applied to BL Lacertae objects and radio jets. Langmuir wave acceleration has also been considered in AGN by Atoyan & Nahapetian (1987). Astrophysical Raman scattering has been detected in the ionosphere through radar scattering experiments (Pineo, Kraft & Briscoe 1960a, b; Salpeter 1960).

In Section 2 the physics of the acceleration and loss mechanisms will be reviewed; although they are well-known to many plasma physicists they were unexploited in the context of astrophysics until the work of Krishan (1983). We apply these processes to regions of varying density in Section 3, generalizing and expanding upon our earlier work (Krishan 1985; Krishan & Wiita 1986). In Section 4 we fit our model to 'average' spectra of Seyfert 1s and quasars and show that they not only provide good matches, but that the values for plasma density and temperature, and the Lorentz factors needed in our model, do not contradict properties of AGN commonly agreed upon. Our conclusions and proposals for additional work are summarized in Section 5. This theory has been applied to the radio spectrum of AGN (Krishan 1988a), and in an alternative explanation of superluminal motion invoking Raman scattering (Krishan 1988b).

2 PHYSICS OF STIMULATED RAMAN SCATTERING

2.1 Fundamental requirements

In our picture an AGN continuum is dominated by high-powered coherent radiation produced by the scattering of a pump field off a collective mode of a relativistic electron beam. This process is known as stimulated Raman scattering (SRS). The pump can be either an electromagnetic wave or a static periodic electric or magnetic field. In more conventional models of AGN, synchrotron photons are analogous to the pump and they undergo inverse Compton scattering to produce high-frequency radiation. But in SRS, the pump field, which in this case can be either cyclotron or synchrotron photons emitted very close to a SMBH or conceivably even thermal photons, for example from an accretion disc, is scattered by a Langmuir wave. The Langmuir wave is a collective excitation of the high-energy electrons, and, therefore, SRS is a faster and more efficient process than inverse Compton scattering. This important point will be discussed in more detail in Section 2.3.

Aside from a source of relatively low-energy photons, most likely of cyclotron origin to provide monochromaticity best, our mechanism works most easily if an initially moderately relativistic electron beam is already present. Observational evidence which is most simply explained if the beams

remain relativistic to large distances from the actual nucleus is growing (Bridle 1988; Barthel 1989). The self-consistent generation of such a beam is beyond the scope of the present paper, although some remarks on this point are made in Appendix A, and we note that such beams are common features of many accretion models for AGN (e.g. Lovelace 1976; Blandford & Payne 1982; Wiita, Kapahi & Saikia 1982; Begelman, Blandford & Rees 1984; Wiita 1985, 1990), and thus think it fair to invoke their existence in this case. We should point out that the entire mechanism could work even in the initial absence of such a beam as electrons could be pulled directly from the tail of the thermal distribution. Since this tail is collisionless, the electrons are easily accelerated by an electric field. Our picture also requires an ambient plasma whose density is somewhat higher than that typically assumed, but it remains significantly below plausible accretion disc densities in the region where such discs are likely to exist ($r \lesssim 10^3 R_s$, with $R_s = 2 GM_{\text{SMBH}}/c^2$), and may serve as the confining medium for the broad emission line clouds.

2.2 Acceleration

2.2.1 Production of Langmuir waves

The Langmuir waves necessary to accelerate the electrons can be produced through Raman forward-scattering. RFS is characterized by the beating of two electromagnetic waves of angular frequency and wave vector $(\omega_0, \boldsymbol{\kappa}_0)$ and $(\omega_1, \boldsymbol{\kappa}_1)$ respectively, so that

$$\omega_0 - \omega_1 = \omega_p, \quad \boldsymbol{\kappa}_0 - \boldsymbol{\kappa}_1 = \boldsymbol{\kappa}_p, \quad (2.1)$$

where $\boldsymbol{\kappa}_p$ is the wave-vector of the longitudinal wave and the plasma frequency, $\omega_p = (4\pi e^2 n/m_e)^{1/2}$, with n the number density of the ambient plasma. The collective oscillations are excited by the ponderomotive force generated by the beating of the two electromagnetic waves. For RFS these waves could possibly grow from thermal noise or synchrotron emission, but would tend to have too broad a bandwidth to generate strong enough Langmuir waves. The most promising type of seed waves in the context of AGN are cyclotron emission, as discussed in Appendix A, where it is shown that there is enough luminosity in the cyclotron lines to power the Langmuir waves through this non-thermal process. Another way to produce Langmuir waves would be through a cyclotron maser, but this would require a severe population inversion; only a mild one is needed (see Appendix A).

The situation of interest has $\omega_0 \gg \omega_p$, so that the phase velocity of the Langmuir wave is

$$v_\phi = \omega_p/\kappa_p = (\omega_0 - \omega_1)/(\kappa_0 - \kappa_1) \approx c(1 - \omega_p^2/\omega_0^2)^{1/2}. \quad (2.2)$$

Such a wave will grow in strength until its amplitude becomes 'relativistic'; in other words, the quiver velocity of the electrons

$$v_Q = eE_L/(m_e \omega_p), \quad (2.3)$$

approaches c (E_L is the induced electric field amplitude). Then the wave traps electrons in the tail of the distribution and accelerates them. Because $v_\phi \approx c$, a sizeable fraction of the electrons will stay in phase with the wave for extended periods and can thus be accelerated to high values of γ . RFS

has been considered in the context of achieving extremely high energies in particle accelerators and experimental evidence for the efficiency of this process is available (Tajima & Dawson 1979; Joshi *et al.* 1981; Tajima 1982; Joshi 1982). The magnitude of the electric field produced through RFS can be shown to be derivable from

$$\nabla \cdot \mathbf{E}_L = -4\pi en', \quad (2.4)$$

where the density fluctuation of electrons bunched by the wave is $n' \lesssim n$ (Tajima 1982). This has the solution

$$|\mathbf{E}_L| = 4\pi en'/\kappa_p = 4\pi en' v_\phi/\omega_p \approx 4\pi en' c/\omega_p. \quad (2.5)$$

We thus generalize the result of Krishan (1985) to

$$E_L = m\omega_p c \xi/e, \quad (2.6)$$

with $\xi = n'/n \lesssim 1$.

Since the rate at which an electron gains energy in a field of strength E_L is

$$dW/dt = e\mathbf{E}_L \cdot \mathbf{v}, \quad (2.7)$$

the time needed to accelerate an electron to reach an energy γmc^2 is approximated by combining equations (2.6) and (2.7):

$$t_a \approx \gamma mc^2/(E_L ce) = \gamma/(\omega_p \xi). \quad (2.8)$$

In a medium whose density changes, a necessary condition for this acceleration to work is that $t_a \ll t_{\text{cross}}$, the time during which a particle moves through a significant change in density

$$t_{\text{cross}} = n/(c|\nabla n|). \quad (2.9)$$

2.2.2 Effectively spatially periodic magnetic fields

The high-amplitude Langmuir waves produced by RFS will be subject to magnetic modulation instabilities (Bel'kov & Tsytoich 1979, 1982; Kono, Skoric & ter Haar 1981; Tsytoich 1983). This is because random small local magnetic fields have as their main effect the production of local changes in the phase of the oscillations. These plasma oscillations with inhomogeneous phase distributions produce vortical currents which tend to increase the spontaneously produced magnetic fields (Bel'kov & Tsytoich 1979; Kono *et al.* 1981). The vortical current is given by

$$\mathbf{J} = -ie(16\pi m\omega_p)^{-1}[\nabla \times (\mathbf{E}_L \times \mathbf{E}_L^*)], \quad (2.10)$$

and the result which is applicable under the conditions of interest is that this magnetic modulation instability excites magnetic fields with spatial periods peaking at $\kappa_i = \omega_p/c$, but ranging up to $\kappa_{\text{th}} = \omega_p/v_{\text{th}}$, with v_{th} the average thermal velocity of the electrons. The resultant field (at κ_i) has an amplitude (Bel'kov & Tsytoich 1979)

$$B_{\text{MM}} = eE_L^2/(4mc\omega_p). \quad (2.11)$$

This quasi-stationary spatially periodic magnetic field will remain excited as long as the Langmuir electric field continues to be reinforced via the beating of two electromagnetic waves.

2.3 Radiation losses

The discussion in Section 2.2 has left us with an ambient plasma of density n , fluctuations of magnitude ξ in that den-

sity due to Langmuir wave bunching, a relativistic beam of electrons of density n_b , accelerated by RFS off the Langmuir waves, and a quasi-stationary spatially periodic magnetic field of amplitude B_{MM} .

2.3.1 Raman back-scattering

Under these conditions the beam electrons will lose energy most rapidly either through stimulated Compton scattering (SCS) or stimulated Raman scattering (e.g. Kaw *et al.* 1976; Hasegawa 1978; Barr, Boyd & Coutts 1988). If the individual fast electrons scatter off photons we have SCS, but if the scattering is by the collective plasma oscillation of the electrons we have stimulated Raman back-scattering (RBS). (The process of scattering off collective oscillations of the ions is called stimulated Brillouin scattering, but is negligible for astrophysical plasmas.) The dominance of RBS over SCS is fundamentally due to the fact that the former goes as $B_{MM}^{2/3}$, while the latter goes as the square of the amplitude, B_{MM}^2 . For the conditions envisaged here, the former is a much faster loss mechanism (Krishan 1983). The analysis is simplest if carried out in the frame in which the beam is at rest; in that frame the beam becomes a stationary plasma and the spatially periodic magnetic field produced by collective effects acts as an electromagnetic wave of frequency $\omega_i = \gamma\kappa_i v_b$ and amplitude $E_i = \gamma(v_b/c)B_{MM}$, with $v_b \approx c$ the beam velocity.

The conditions for RBS are

$$\omega_i - \omega_s = \omega_b \quad \text{and} \quad \kappa_i - \kappa_s = \kappa_b, \quad (2.12)$$

where ω_s is the angular frequency of the scattered wave and ω_b is the plasma frequency of the beam electrons, all in the beam frame. The collective nature of the scattering is lost for very large values of $\kappa_b \approx \kappa_D \equiv 2\pi/\lambda_D$ (with λ_D the Debye wavelength); this is the condition of heavy Landau damping, and implies that SCS takes over from RBS. After computing the dispersion relation in the beam frame and then transforming back into the laboratory frame, the rate of loss of beam electrons to photons via RBS can be shown to be (Hasegawa 1978; Krishan 1983)

$$\Gamma_L = \gamma^{-1} (2\beta_i^2 \omega_b^2 \gamma \omega_p)^{1/3}, \quad (2.13)$$

where $\beta_i \equiv eB_{MM}/(m\omega_p c)$, as long as $\beta_i^2 \gg \omega_b/\gamma\omega_p$. Using equations (2.6) and (2.11) we find

$$\beta_i = eE_L^2/(4m^2 c^2 \omega_p^2) = \xi^2/4. \quad (2.14)$$

Any ambient large-scale magnetic field, such as that generating cyclotron or synchrotron seed photons, could act as a guide to the beam, but does not affect this scattering process.

2.3.2 Emitted power

We now introduce the key assumption that the rate at which beam electrons lose energy via RBS, (2.13), equals the rate at which they gain energy from RFS, obtainable from (2.8). With this very reasonable assumption, $\Gamma_L = \Gamma_G \equiv t_a^{-1}$, and a little algebra, we can actually relate the density of the electron beam to the ambient density and the size of density fluctuations:

$$n_b/n = 8/(\xi\gamma). \quad (2.15)$$

Note that, while the ambient plasma has an important role in

the acceleration, only the relativistic beam particles are responsible for the emitted radiation, so that, in general, $n_b < n'$. Further, this ansatz has the pleasant consequences that a steady-state distribution is produced, and the spread in γ at a given density is small. This is helpful in reducing assorted plasma instabilities.

The total power emitted by RBS can be closely estimated by noting that, as the back-scattering power rises, the effective electric field ($\beta_i \times \mathbf{B}_s$) produced by the scattering also increases. This field will trap the beam electrons and increase their thermal spread. It is simple to show that the effective thermal spread engendered by this trapping electric field is

$$v_{TRBS} = [2e|\beta_i||B_s|/(\kappa_b m)]^{1/2}, \quad (2.16)$$

where κ_b is the wave vector of the longitudinal wave excited in the beam electrons responsible for scattering the photons in RBS, and B_s is the scattered field (in the beam frame). The growth in the field is halted and the power must saturate when $\kappa_b \approx \omega_b/v_{TRBS}$; this is because at that point the plasma wave (ω_b, κ_b) suffers heavy Landau damping. As mentioned previously, this implies that the radiation process shifts from scattering of the spatially periodic magnetic field by the plasma wave (RBS) to scattering of that field by individual electrons in the beam (inverse Compton scattering) (Hasegawa 1978; Krishan 1985). Hence, the maximum value of B_s is obtained from the conditions

$$\kappa_b = \kappa_D = \omega_b \gamma / (c\Delta\gamma), \quad (2.17)$$

since $c\Delta\gamma/\gamma$ is the thermal spread of the beam electrons in the beam frame. Using (2.16) we find

$$B_s = m\omega_b^2 / (2e\beta_i \kappa_b). \quad (2.18)$$

It is then easily shown that the total power emitted in the 'laboratory' frame, where the beam has velocity $\approx c$, is (Krishan 1985)

$$L = (c/4\pi) 4B_s^2 \gamma^2 A = (\Delta\gamma/\gamma)^2 mc^3 \gamma^2 \beta_i^{-2} n_b A, \quad (2.19)$$

with A the cross-sectional area of the beam (or beams). This energy is emitted essentially at the RBS scattered wave frequency,

$$\omega_s = 2\gamma^2 \kappa_0 c = 2\gamma^2 \omega_p, \quad (2.20)$$

for, just as in the standard derivation of inverse Compton scattering, one power of γ arises in each of the Lorentz transformations into and out of the beam frame (Hasegawa 1978). Upon equating ω_s to the maximum frequency of emission found from (2.17), we have $\omega_s = 2\gamma^2 \omega_p \approx \omega_{cr} = \gamma^2 \omega_b/\Delta\gamma$, or

$$(\Delta\gamma/\gamma)^2 \approx \omega_b^2 / (4\gamma^2 \omega_p^2) = 2/(\gamma^3 \xi), \quad (2.21)$$

where the last step follows from (2.15) and the definition of plasma frequency. Substituting (2.21) into (2.19) and using (2.14), we obtain the final expression for the emitted power at $\omega = \omega_s$:

$$L = 256 mc^3 n A \gamma^{-2} \xi^{-6} = 6.28 \times 10^6 n A \gamma^{-2} \xi^{-6} \text{ erg s}^{-1}, \quad (2.22)$$

the last equality holding if n and A are in cgs units.

3 GENERAL PROPERTIES OF THE EMISSION

Thus the key result of our proposed mechanism is that denser regions of the plasma, presumably closer to the very

core of the AGN, produce higher frequency SRS emission. This is because the frequency of emission scales with ω_p , or $n^{1/2}$, if γ^2 does not vary very much. Of course the total power also depends on A , the cross-section for beam-plasma interaction, and ξ , the fraction of bunched electrons. It is interesting to note that the power can be made independent of frequency (corresponding to $\alpha = 1$, close to the observed α_{IX}), if γ , ξ and the product nA are constant. If the beams have a constant opening angle so that $A \propto r^2$, then this special case corresponds to $n \propto r^{-2}$, the familiar law of density variation in a uniformly expanding or contracting medium.

Of course in reality we do not expect such neat constancy, but if this process operates as we have proposed, then ξ should not vary too much and ought to be within an order of magnitude of unity. Further, the evidence that typical observed radio jets have fairly constant opening angles over large ranges is quite good (Spangler 1979; Kerr *et al.* 1981; Gopal-Krishna & Wiita 1987). Thus, if we write $A \equiv \Omega r^2$, it is not unreasonable to simplify the analysis by choosing $\Omega = \text{constant}$; noting that a factor of 8π has already been pulled out of the total surface area at radius r if two beams are present, $0.03 \leq \Omega \leq 1$ is likely.

Following Krishan (1985) we then parameterize the spatial variations in density and Lorentz factor as

$$n(r) = n_i (r_i/r)^{q+2}, \quad (3.1a)$$

$$\gamma^2(r) = \gamma_i^2 (r_i/r)^p, \quad (3.1b)$$

where r_i is some fiducial value and n_i and γ_i are the density and Lorentz factor at that point. It will later prove convenient to take r_i at the position that produces radiation whose frequency corresponds to the break between α_{IX} and α_X . Inserting (3.1) into (2.20) and (2.22) we find

$$\omega(r) = \omega_i (r_i/r)^{p+1+q/2}, \quad (3.2a)$$

$$L(r) = (r_i/r)^{q-p}, \quad (3.2b)$$

and

$$f_\nu(r) \equiv L(r)/\omega(r) = (L_i/\omega_i)(\omega/\omega_i)^{-\alpha}, \quad (3.2c)$$

where

$$\omega_i = 2 C_1 \gamma_i^2 n_i^{1/2}, \quad L_i = C_2 \Omega \xi^{-6} n_i r_i^2 \gamma_i^{-2},$$

$$C_1 = (4\pi e^2/m_e)^{1/2} = 5.64 \times 10^4, \quad C_2 = 256 mc^2 = 6.28 \times 10^6,$$

and

$$\alpha = (2p+1-q/2)/(p+1+q/2), \quad (3.2d)$$

so that the special case mentioned above had $p = q = 0$.

At this point we note that, because the simple form of the loss rate Γ_L was obtained subject to the approximation $\beta_i^2 \gg \omega_b/(\gamma\omega_p)$, we find that a self-consistency relation between ξ and γ follows from (2.14) and (2.15)

$$\xi > 2^{2/9} (8/\gamma)^{1/3}. \quad (3.3)$$

Examination of equations (3.2c and d) shows that for appropriate choices of p and q a single constant spectral index is produced. Therefore the observed breaks in the AGN spectrum require either a modification of at least one of these parameters or the intervention of additional physical processes not yet considered. As will be discussed in Section 4 in some detail, the key break at sub-keV energies apparently must arise from a change in the density law or the

variation of γ with distance (or an unmodelled variation in Ω or ξ). Such changes are related to the magnetohydrodynamics of the in- and out-flows in the neighbourhood of the compact object (e.g. Rees 1984). This is an extensive problem in itself and is beyond the scope of this paper. Of course, similar lacunae are present in essentially all previous explanations of the AGN continuum.

On the other hand, the breaks at both very high- and very low-energies, necessary to avoid over-producing γ -rays above a few MeV (e.g. Swanenburg *et al.* 1978) and to yield the usual flat radio spectra, emerge quite naturally. The constraints on the sizes of emission regions which emerge from observations (see Sections 4.1 and 4.2), combined with the necessity of keeping the density from rising unrealistically rapidly upon approach towards the SMBH, indicate that γ -ray production should become much weaker above ~ 2 MeV. Further, pair production and other cooling mechanisms (Zdziarski 1986; Begelman 1988) will also act to produce a significant break in the spectrum in that range.

While a formal extrapolation of equations (3.2) down to frequencies below 10^{12} Hz would imply a steep spectrum continuing into the radio, the normally invoked process of synchrotron self-absorption will come into play and can significantly cut the net emission in the radio. The fact that our density distribution is naturally falling off with distance allows the flatter spectra characteristic of inhomogeneous self-absorption (e.g. Marscher 1977) to be consistently applied. Additional modifications to the soft end of the spectrum may be produced by anomalous absorption (Krishan 1987) or by Raman scattering in the accretion disc (Krishan 1988a).

The polarization of the emitted radiation in the model corresponds to that of the scattering magnetic field which is generated by the magnetic modulation instability, as mentioned in Section 2.2.2. The polarization of the field is in general elliptical, and is produced from the current of equation (2.10). However, we do not expect any large net polarization, as the actual Langmuir field, as well as the generated magnetic field, will be the superposition of many such small-scale fields, whose polarizations would tend to cancel. Thus, despite the coherent nature of the processes envisaged here, they are likely to produce the typically low polarizations observed in AGN (e.g. Stein & O'Dell 1985). None the less, significant changes in polarization are possible if, during a fluctuation, a particular sub-area temporarily dominates the emission at a particular frequency.

Variability of the intensity of emission would naturally be produced through the density fluctuations to be expected in magnetohydrodynamical flows. The length scale of such strong density fluctuations (with $\delta n \sim n$) vary from $\sim R/3$ all the way down to the collisional mean-free-path of a proton, where R is the radius of the plasma region. The time-scales of such fluctuations thus vary from $\sim R/v \approx R/c$ down to the proton-proton collision time, assuming the fluctuations are coherent over the limited range of radii from which emissions at a given frequency emerges. Therefore, we would expect to observe time variability over a rather broad range of period at a given frequency. Since the value of R rises as the frequency falls, the most rapid variability is expected in the X-ray or γ -ray band. Because the collisional frequency is also highest at higher densities, this mechanism permits extremely rapid fluctuations at higher frequencies ($\lesssim 1$ s is

possible in X-rays) while coherent optical variability on time-scales of a few minutes (despite crossing times typically light-days or more for the optical emission zone, see Section 4.2) is possible. In general, the range of times (on a logarithmic scale) over which significant fluctuations are expected will shrink with decreasing frequency.

4 COMPARISONS WITH OBSERVATIONS

4.1 General constraints

If a general model for AGN exists at all, it involves a central engine, presumably a SMBH of between $\sim 10^6$ and $\sim 10^9 M_\odot$, producing a maximum radiative output of $\leq L_{\text{Edd}}$, where $L_{\text{Edd}} \approx 1.3 \times 10^{46} M_8 \text{ erg s}^{-1}$, and M_8 is the mass of the SMBH in units of $10^8 M_\odot$. Broad-line emission is usually taken to imply the existence of many rather small, mostly ionized clouds scattered in a region ranging from ~ 0.01 to ~ 1 pc with temperatures $\sim 10^4$ K and densities about $\sim 10^8$ to $\sim 10^{11} \text{ cm}^{-3}$. These BLR clouds, providing a covering factor of ~ 0.1 , are presumably confined by a hotter, but less-dense medium, and the clouds are almost certainly photoionized by the hard UV and X-ray emission from a small region surrounding the central engine (e.g. Mathews & Capriotti 1985; Ferland & Shields 1985; Wiita 1985; Krolik 1988).

Because our model predicts that the highest frequency emission must emerge from the densest regions closest to the central engine, the most fundamental constraints arrived at from the previous paragraph can be written as

$$R_S = 3.0 \times 10^{13} M_8 \text{ cm} > 2.3 \times 10^{13} L_{46} \text{ cm}, \quad (4.1)$$

$$R_S \leq r_\gamma < r_X < r_{\text{UV}} < r_o < r_{\text{IR}}, \quad (4.2)$$

$$r_{\text{UV}} < R_{\text{BLR}} \text{ and} \quad (4.3)$$

$$n(R_{\text{BLR}}) \ll n_{\text{BLR}}, \quad (4.4)$$

where L_{46} is the total power in units of $10^{46} \text{ erg s}^{-1}$, r_{band} is the average radius at which that type of emission is generated in the SRS picture ('o' standing for optical), R_{BLR} is the typical distance of the broad line clouds from the centre, n_{BLR} is the average density in such clouds, and $n(R_{\text{BLR}})$ is the density implied by our model at R_{BLR} .

For our model to be viable, other sources of emission, particularly bremsstrahlung, should not dominate our proposed continuum, and we must ascertain that the radiation produced at r_{band} can reach observers outside the emitting galaxy (except for the fraction intercepted by the BL clouds, and at greater distances, the narrow-line clouds). Further, our models should involve plausible accretion rates and the self-consistency relations given in equations (2.8), (2.9) and (3.3) must be satisfied.

4.2 Fits to 'average' AGN spectra

4.2.1 Quasars

Although, as discussed in Section 1, quasars exhibit a significant range in continuum properties, especially in the X-ray band, we still feel it useful to attempt to describe an over-all 'average' spectrum, and see if the SRS picture can plausibly satisfy the constraints mentioned in Section 4.1. Tananbaum *et al.* (1983) provide a set of uniform observations of 3CR quasars at X-ray (0.5–4.5 keV), optical (2500 Å) and radio

(5 GHz) frequencies. From this compilation we find an average value for $L_x \approx 6 \times 10^{45} \text{ erg s}^{-1}$ at an average frequency of $\nu_x = 4.8 \times 10^{17} \text{ Hz}$ (2 keV, $\omega = 3.0 \times 10^{18} \text{ rad s}^{-1}$). Using the 'typical' soft X-ray spectrum of $\alpha_x = 0.7$ (e.g. Petre *et al.* 1984), we extrapolate down to a band centred on 0.2 keV, where we assume that the steeper IR-X and flatter X-ray continua join, and find $\langle L_{0.2 \text{ keV}} \rangle \equiv L_i \approx 3 \times 10^{45} \text{ erg s}^{-1}$. Using $\alpha_{\text{IX}} = 1.2$, we find the continuum $\langle L_o \rangle = 6.9 \times 10^{45} \text{ erg s}^{-1}$ (at 4000 Å), although the average total optical emission from Tananbaum *et al.* is about $1.9 \times 10^{46} \text{ erg s}^{-1}$. We adopt the former value and attribute the significant excess to the average of the 'big blue bump', for this difference is completely consistent with the measurements of Malkan & Sargent (1982) and Malkan (1983), when one notes that a positive correlation between the amount of the quasi-thermal blue bump spectrum and the quasar luminosity appears to be present. Despite the actual wide variation in the hard X-ray spectra of quasars, for concreteness, and to provide a stringent test of our model, we have assumed that $\alpha_x = 0.7$ holds out to 2 MeV, giving $\langle L_\gamma \rangle \approx 4.7 \times 10^{46} \text{ erg s}^{-1}$, $\langle L_{\text{tot}} \rangle \approx 10^{47} \text{ erg s}^{-1}$, and implying $M_{\text{BH}} \approx 10^9 M_\odot$ and $r_\gamma \geq 3 \times 10^{14} \text{ cm}$ from (4.1).

For such a powerful object the BLR would probably range between 0.1 and 1 pc, and a reasonable radius at which optical emission could be produced would be $r_o = 10^{17} \text{ cm}$, since then the UV and X-ray emission will emerge from well inside the BL clouds (*cf.* Section 4.3). We have found that choosing the break radius at $r_i = 2 \times 10^{16} \text{ cm}$ yields a reasonable solution, although the results we shall now give are not unique. Using the relations $L_i = C_2 \Omega \xi^{-6} n_i r_i^2 \gamma_i^{-2}$ and $\omega_i = 2 C_1 \gamma_i^2 n_i^{1/2}$, we can find n_i and γ_i in terms of r_i , L_i and ω_i

$$n_i = (\omega_i / 2 C_1 C_2)^{2/3} R_i^{-4/3} \Omega^{-2/3} \xi^4 L_i^{2/3} \quad (4.5)$$

$$= 2.64 \times 10^{12} R_{16}^{-4/3} \Omega^{-2/3} \xi^4 L_{45}^{2/3} (\omega_i / 3.0 \times 10^{17})^{2/3} \text{ cm}^{-3};$$

$$\gamma_i = C_2^{1/6} / (2 C_1)^{1/3} R_i^{1/3} \Omega^{1/6} \xi^{-1} L_i^{-1/6} \omega_i^{1/3} \quad (4.6)$$

$$= 1.29 \times 10^3 R_{16}^{1/3} \Omega^{1/6} \xi^{-1} L_{45}^{-1/6} (\omega_i / 3.0 \times 10^{17})^{1/3},$$

where $R_{16} = R_i / 10^{16} \text{ cm}$ and $L_{45} = L_i / 10^{45} \text{ erg s}^{-1}$. In order to determine reasonable values for Ω and ξ (or actually, their combination $\Omega(\xi^6)$) we impose the constraint (4.4) by choosing $n_o = 10^9 \text{ cm}^{-3}$, so $\gamma_o = 1.15 \times 10^3$. Putting $L_o = 6.9 \times 10^{45} \text{ erg s}^{-1}$ into (2.22) yields $\Omega / \xi^6 = 145$. For concreteness, let $\Omega = 1$, so that $\xi = 0.436$, and the constraint (3.3) is satisfied but even $\Omega = 0.03$ allows an acceptable value of ξ . Inserting $R_{16} = 2$, $L_{45} = 3$ and $\Omega / \xi^6 = 145$ into equations (4.5) and (4.6) yields $n_i = 7.9 \times 10^{10} \text{ cm}^{-3}$ and $\gamma_i = 3.1 \times 10^3$.

We are now able to find the values of p and q which are appropriate outside r_i , by simultaneously solving equations (3.2a) and (3.2d), because ω_o and ω_i as well as α are all input parameters. Alternatively, we can compare n_i with n_o and γ_i with γ_o using equations (3.1a and b). Naturally, both give identical results: $q = 0.714$, $p = 1.232$, or $n \propto r^{-2.714}$ and $\gamma \propto r^{-0.616}$.

The same procedure can be followed in investigating plausible solutions for the production of high-energy photons. Using identical values for r_i , n_i , γ_i and L_i , but with $L_\gamma = 4.7 \times 10^{46} \text{ erg s}^{-1}$, $r_\gamma = 3 \times 10^{14} \text{ cm}$, and $\omega_\gamma = 3.0 \times 10^{21} \text{ rad s}^{-1}$, we find that a steeper rise in density towards the very centre is required, with $q = 1.234$, $p = 0.576$, and $n_\gamma = 6.2 \times 10^{16} \text{ cm}^{-3}$ and $\gamma_\gamma = 1.0 \times 10^4$.

4.2.2 Seyfert 1 galaxies

Many more data are available for Type 1 Seyfert galaxies, particularly in the X-ray band, and it is for them that the 'universal' spectrum is more clearly established (e.g. Zdziarski 1986). We mainly used the compilation of data in Lawrence & Elvis (1982) who give values for X-ray emission in both the 0.5–4.5 and 2–10 keV bands as well as the luminosities in the O III ($\lambda 5007$) and H β ($\lambda 4861$) lines and fluxes at 3.5 and 10 μm . Using these data for soft X-rays and assuming (as before) that the values for α are 1.2 and 0.7 below and above 0.2 keV, we find: $\langle L_x \rangle = 3.0 \times 10^{43} \text{ erg s}^{-1}$, $\langle L_0 \rangle = 3.5 \times 10^{43} \text{ erg s}^{-1}$, $\langle L_i \rangle = 1.5 \times 10^{43} \text{ erg s}^{-1}$ and $\langle L_\gamma \rangle = 2.4 \times 10^{44} \text{ erg s}^{-1}$. These numbers are in decent agreement with the data given in Lawrence & Elvis at other frequencies, but the lack of information on observational bandwidths in that paper or in the bulk of the references cited by them precluded our obtaining more accurate broad-band luminosity values from the quoted narrow-band luminosity or flux values.

Since these Seyferts are about two orders of magnitude weaker than the quasars considered above, the central SMBH should be less massive and all relevant radii would also shrink, with reasonable values for $R_s \approx 3 \times 10^{12} \text{ cm}$, $R_{\text{BLR}} \approx 5 \times 10^{16} \text{ cm}$. This implies that choosing $r_0 = 10^{16} \text{ cm}$, $n_0 = 10^9 \text{ cm}^{-3}$ and $r_i = 10^{15} \text{ cm}$ gives us a good chance to fit the spectra subject to the constraints (4.1–4.4). Following the procedure of the previous subsection, we find the following set of parameters can yield a decent fit

$$r_i = 10^{15} \text{ cm}, L_i = 1.5 \times 10^{43} \text{ erg s}^{-1}; n_i = 2.0 \times 10^{11} \text{ cm}^{-3},$$

$$\gamma_i = 2.5 \times 10^3;$$

$$r_0 = 10^{16} \text{ cm}, L_0 = 3.5 \times 10^{43} \text{ erg s}^{-1}; n_0 = 1.0 \times 10^9 \text{ cm}^{-3},$$

$$\gamma_0 = 1.1 \times 10^3;$$

$$r_\gamma = 10^{13} \text{ cm}, L_\gamma = 2.4 \times 10^{44} \text{ erg s}^{-1}; n_\gamma = 2.7 \times 10^{17} \text{ cm}^{-3},$$

$$\gamma_\gamma = 7.2 \times 10^3;$$

$q = 0.299$, $p = 0.660$ for $\omega < \omega_i$ while $q = 1.067$ and $p = 0.467$ for $\omega > \omega_i$; in this case $\Omega \xi^{-6} = 73$, so, for $\Omega = 1$, $\xi = 0.489$. Therefore, just as for the quasar case, to obtain the break in the spectrum in the sub-keV range we must have the density law steepen in the innermost region, if we assume ξ and Ω are constant throughout, which is of course an oversimplification. None the less, the values of p and q given above for QSOs and those given here for Seyfert 1s illustrate the types of density and Lorentz factor dependences that ought to be present in the region surrounding SMBHs and are at least as reasonable as those used in other models for AGN continua. The values required for γ , typically 10^3 – 10^4 , are similar to those obtained in other calculations (e.g. Zdziarski 1986).

4.3 Consistency checks

The generation of Langmuir waves, the production of periodic magnetic fields and the radiation process all occur on the extremely short time-scales given by equations (2.8) and (2.13). Therefore, much slower processes such as synchrotron and inverse Compton emission would be insignificant compared to the SRS emission, except perhaps at the lowest (radio) frequencies.

However, it is not immediately apparent that bremsstrahlung emission would also be negligible, and, as a matter of fact, it can be quite significant. But, of course, bremsstrahlung depends upon the temperature of the plasma, which we have so far not had to specify. If the gas in the BLR has $n \leq 10^7 \text{ cm}^{-3}$, and does confine the clouds, pressure balance demands that $T > 10^6 \text{ K}$ in that region, while the assumption that the gas is well ionized everywhere implies that $T \geq 10^4 \text{ K}$ in all regions. If the temperature were to be taken as $\geq 10^6 \text{ K}$ throughout the volume of the plasma, then the total emission in soft X-ray bremsstrahlung would greatly exceed the SRS-produced X-rays and our model would not dominate the spectrum in that band since the entire plasma cloud is optically thin to such X-rays. This alternative is certainly possible, and might conceivably obviate the necessity for a break in the density power law to explain the excess X-ray emission. We plan to examine this possibility in detail in future work.

Still, for the simplest version of our model to hold, we must assume that the temperature decreases as the gas approaches the central object, to that in the densest regions it might be comparable to the typical temperature expected of an accretion disc of a few ten thousands of degrees (e.g. Malkan 1983). The higher temperatures further away might be understood in the context of a corona above the disc (e.g. Begelman 1985). Then, the innermost regions of highest density would nominally produce extremely large amounts of optical (actually soft UV) bremsstrahlung, but because our entire plasma distribution is optically thick to such emission, very little of it escapes and the bulk of it merely serves to thermalize the plasma. The higher frequency bremsstrahlung emission produced by the hotter gas at larger radii is quite small in comparison to the SRS emission produced closer to the central engine, so that the transparency of the plasma at these frequencies only means that our SRS-produced photons can escape directly.

Our model thus includes a new possible explanation for the big blue bump which must differ from that of Malkan (1983) in that the optical and UV quasi-thermally distributed photons produced in the inner parts of an accretion disc would not be able to emerge through plasma densities of the type envisaged here. Preliminary calculations indicate that the small fraction of the total optical and UV bremsstrahlung that does escape from the inner cloud is of the appropriate magnitude and spectrum to account for the excess above the power law often seen in quasars. It is less important in the case of Seyferts, which is also in accord with observations. However, more detailed computations will be required before we can make definite predictions as to the detailed nature of this excess flux, particularly in view of the complexities added by the forest of Fe II lines in the BLR (e.g. Wills, Netzer & Wills 1985). We can finally mention in this regard that the soft X-ray bremsstrahlung could be important in sources where there are indications of large soft X-ray and EUV excesses (e.g. Bechtold *et al.* 1987); in our picture the gas temperature would remain $\sim 10^6 \text{ K}$.

Another point which should be checked involves the total mass required to be in the form of our plasma, for the higher than usually considered densities might lead to unacceptably high values. However, total gas masses of between $\sim 10^3$ and $\sim 10^4 M_\odot$ out to distances of 10^{19} cm are implied by our density distributions for Seyferts and quasars. Such values,

while much greater than the mass necessarily involved in the emission line clouds (Mathews & Capriotti 1985), are very small when compared to either the mass of the central black hole or the mass of stars expected to reside within a few parsecs of that monster; we know of no reason why they should not be acceptable.

The extremely efficient nature of the acceleration process assumed in this paper also means that, despite fairly high values of γ and rather rapid variations in density, the acceleration time t_a (equation 2.8) is always many orders of magnitude less (about 10–12) than the density zone crossing time (equation 2.9), so a necessary condition for our mechanism is preserved. Our model does not place any constraint on the accretion rate other than the usual one that the ultimate source of energy should be the infall of some mass into the SMBH, so that for an efficiency of ϵ , $\dot{M} = L\epsilon^{-1}c^{-2}$; as in standard accretion pictures, $\epsilon \sim 0.1$ is expected, so for $L_{46} = 1$, $\dot{M} \approx 2 M_{\odot} \text{ yr}^{-1}$.

4.4 The efficiency of RFS

The theory of RFS employed in Section 2.2 assumes coherence of the two electromagnetic waves beating to produce the Langmuir wave. There are essentially two ways in which these waves could be incoherent. First, the incoherence could be inherent in the generating process; for example, if the waves are produced through cyclotron or synchrotron emission the pulse width ($\Delta t \sim 1/\Delta\omega$) is a source of incoherence. Thermal generation will certainly introduce larger incoherence. A coherent field can also become incoherent while propagating through a turbulent and inhomogeneous medium. Complete discussions of these points are major projects in and of themselves, and we plan to address them in future work; at this point we wish to quantify the allowed degree of incoherence.

The effect of incoherence of the electromagnetic waves is to increase the threshold and reduce the growth rate of RFS. Parametric excitation of waves with stochastic phases has already been studied, mainly to find the average values of the phase amplitudes. It is generally found that the results are modified only when the coherence time becomes shorter than the interaction time. This is equivalent to noting that the power of the pump is spread over a frequency range $\Delta\omega$, while only that within the resonance, of width Γ , is available to drive the instability so that the growth rate is reduced by a factor $\Gamma/\Delta\omega$ (Thomson 1975). Therefore the theory of RFS as used here is valid when $\Delta\omega \lesssim \Gamma$. The growth rate of RFS is given by (Joshi *et al.* 1984)

$$\Gamma = (\omega_p/4c^2)[v_0(0) \cdot v_0(1)], \quad (4.7)$$

where $v_0(0) \approx v_0(1) = eE_0/(m\omega_{0,1}c)$, so using (2.1) our condition is

$$\Delta\omega/\omega_1 \approx \Delta\omega/\omega_p \lesssim (v_0/2c)^2. \quad (4.8)$$

The amplitude of the plasma wave E_L is (Joshi *et al.* 1984)

$$eE_L/(m\omega_p c) = (16v_0^2/3c^2)^{1/3} = \xi, \quad (4.9)$$

where we have used (2.8) to obtain the second equality. Now

the constraint on the bandwidth (4.8) becomes

$$\Delta\omega_{\text{max}}/\omega_p \lesssim 3\xi^3/64 \approx 5 \times 10^{-3}, \quad (4.10)$$

for the values of ξ obtained in Section 4.2. Since the effective bandwidth of this exciting synchrotron radiation can be approximated by $\Delta\omega/\omega \approx 1/\gamma < 10^{-3}$, we may be justified in using synchrotron emission to induce RFS in the coherent approximation. The monochromaticity of cyclotron lines under the conditions outlined in Appendix A satisfies this constraint even more easily.

This limit allows some scope for partial incoherence which may be induced by turbulence. Furthermore, it has been shown that, using the random phase approximation, a finite coherence length can actually restore the growth of the instability in the presences of low-frequency turbulence which is probably inherent in AGN conditions (Laval, Pellat & Pesme 1976; also *cf.* Barr *et al.* 1988).

We must also require that the acceleration time of particles (2.8) is greater than the growth time of the RFS Langmuir waves (4.7), i.e., $t_a = \gamma/(\omega_p \xi) > \Gamma^{-1}$. Substituting for ξ from (4.9) and for Γ from (4.7) we obtain $(v_0/c)^2 > 18.5/\gamma^{3/2}$. Our result of $\xi = 0.45$ ensures that the above condition is met, so there is sufficient time to accelerate the electrons. Further, for this value of ξ one gets $(v_0/c)^2 = 1.7 \times 10^{-2}$. Now using the relations $v_0/c = eE_0/(m\omega_0 c)$ and $L_{\text{syn}}/(4\pi R^2 c) \approx E_0^2/4\pi$, we find that for values of $\omega_0 \gtrsim 2\omega_p$ in high-density regions of our models we get L_{syn} comparable to the total average luminosity of an AGN. In other words, synchrotron photons can indeed provide all necessary seed photons for RFS, and can be provided through typical accretion disc models. Similar statements hold if the underlying energy source first emerges in the form of cyclotron emission. Although this last result is very rough, it is encouraging, and we stress that the overall model appears to satisfy all self-consistency requirements.

5 DISCUSSION AND CONCLUSIONS

Although we believe that we have established the viability of the stimulated Raman scattering mechanism with respect to the production of the overall continuum in active galactic nuclei, we also recognize that much more work will be necessary before it is likely that this process will be accepted as correct. As of now we do not have a natural explanation of why the change in density law should always, or at least usually, occur at a radius corresponding to sub-keV photons. Detailed calculations of the temperature profile, and thus the bremsstrahlung emission, are also desirable, as are quantitative results concerning polarization and variability. The self-consistent generation of adequate seed photons for RFS can be produced through cyclotron lines excited in the hot ambient plasma (see Appendix A), but further exploration of the range of parameters for which these processes are efficient is required.

This model has some very nice properties which distinguish it from most other proposed generators of the continuum for AGN. One key point is that the acceleration and radiation processes are simultaneous and complementary; they naturally produce steady-state values for the fast particles' energies, essentially regardless of the injection spectrum. This also means that the same group of particles could give rise to a wide range of photon energies as they

flow outwards into regions of lower density. Relatively minor changes in the density distribution can produce the observed break in the sub-keV region.

Our conclusions can be summarized as follows.

(i) Cyclotron, or possibly synchrotron or thermally generated, photons can beat in order to produce Langmuir waves through Raman forward-scattering. The cyclotron mechanism is most likely in that these lines have enough luminosity to drive the waves and it also ties together the spatial variations of the plasma density and the magnetic field. These Langmuir waves can accelerate electrons to Lorentz factors of $\sim 10^3$ – 10^4 .

(ii) The magnetic modulation instability generates locally spatially periodic magnetic fields which act as a pump (or soft photons) for Raman backward-scattering.

(iii) RBS can provide the bulk of the continuum emission from the far-IR through γ -rays, with higher energy photons emitted at smaller radii.

(iv) Matching the rate of loss of energy via RBS to the rate of gain of energy from the RFS-generated Langmuir waves keeps the electrons continuously accelerated. This means that a steady-state situation is easily produced, and the density of fast electrons is found in terms of the ambient plasma density.

(v) The coherent nature of this process implies that the thermal spread ($\Delta\gamma/\gamma$) will be small, and the beam is, therefore, resistant to many disruptive instabilities.

(vi) A constant spectral index of unity would be produced in a uniformly expanding plasma with no variation in γ . The observed spectra of AGN imply density decreases with radii characterized by power laws in the range 2.2–3.3.

(vii) Synchrotron emission and synchrotron self-absorption, perhaps coupled with anomalous absorption, probably produce the flat radio spectra. The cut-off in the γ -ray spectrum is probably due to pair production, coupled with general relativistic effects in the immediate vicinity of the central engine.

(viii) Escaping bremsstrahlung photons could explain the big blue bump.

(ix) Variability on time-scales shorter than the light-travel time across a particular radius are possible, as coherent MHD modes can affect the local density, and thus the luminosity, at particular frequencies.

(x) Average net polarization should be small, but significant (a few per cent) values are possible for relatively short periods.

(xi) All of the above properties are intriguing enough to justify additional investigations of this emission mechanism in the context of AGN.

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APPENDIX A: CYCLOTRON ORIGIN OF SEED FIELDS

Although the discussion in Section 4.4 indicates that synchrotron emission may be capable of providing the electromagnetic radiation needed to generate Langmuir waves, it is certainly true, as pointed out by the referee, that

the bandwidth of such photons would be much broader than those used in the laboratory experiments referred to in Section 2.2.1. Here we argue that cyclotron radiation due to non-relativistic electrons can provide an alternative source of EM waves with an adequately narrow bandwidth. Defining $\omega_H = eB/mc = 1.76 \times 10^7 B \text{ s}^{-1}$ as the cyclotron frequency in an ambient field B , the power in a line at $\omega = s\omega_H$, with s an integer, is roughly given by (e.g. Zheleznyakov 1970):

$$L(\omega) = 2[(1 - \omega_p^2/\omega^2)^{1/2}]^{2s-1} e^2 \omega_H^2 \beta_p^{2s} s^{2s+1} (s+1) N \times [c(2s+1)!]^{-1}, \quad (\text{A1})$$

with the component of the velocity perpendicular to the field given by $\beta_p c$, subject to $\beta_p \ll 1$, and N the total number of radiating electrons in that region. For there to be Langmuir wave generation from such harmonics we require $\omega_p = j\omega_H$, with j an integer, and then obtain (in cgs units)

$$L(s\omega_H) = (4.61 \times 10^{-21} n) [1 - (j/s)^2]^{2s-1} \beta_p^{2s} s^{2s+1} (s+1) \times N [j^2(2s+1)!]^{-1}. \quad (\text{A2})$$

Note that cyclotron emission at the fundamental frequency is quite likely to be suppressed through absorption by electrons which have the same sense of rotation as that of the wave, and in fact solar data give evidence for the higher harmonics being stronger (Zheleznyakov 1984).

The most extreme parameters occur if the γ -rays are to be produced in this manner. As an illustration, make the assumption that a population of supra-thermal electrons with $\beta_p = 0.3$ exists, and use the values obtained in Section 4.2.2, $n = 2.7 \times 10^{17} \text{ cm}^{-3}$, $r = 10^{13} \text{ cm}$, so $\omega_p = 2.9 \times 10^{13}$, to give an estimate of the required fields. If $j=1$ and $s > 1$, we can, for example, satisfy equation (2.1) with $\omega_0 = 3\omega_p = 3\omega_H$ and $\omega_1 = 2\omega_p = 2\omega_H$, so that (A2) yields

$$L(3\omega_H) = 4.34 \times 10^{-24} nN = 4.34 \times 10^{-24} n^2 V \text{ erg s}^{-1}, \quad (\text{A3})$$

with V the volume producing this radiation. This should be comparable to, but somewhat less than, the total required luminosity of the EM waves which can be estimated as (cf. Section 4.4)

$$L(\omega_0) = cr^2 E_0^2 = c(v_0 mr/e)^2 \omega_0^2 = 2.6 \times 10^{45} \text{ erg s}^{-1}, \quad (\text{A4})$$

for these parameters. Combining (A3) and (A4), we find $V \lesssim 8.2 \times 10^{33} \text{ cm}^3$, which is a minute fraction of $r^3 = 10^{39} \text{ cm}^3$, indicating that the requisite emission can be produced in a narrow shell or segment of a cone over which the magnetic field strength, and thus ω_H , is effectively constant. Equivalently, a smaller density of supra-thermal particles in a larger (but still relatively small) volume may provide the same luminosity. Assuming the same emission volume for the other harmonic leads to $L(2\omega_H)$ within a factor of two of $L(3\omega_H)$, so that the combination of these photons to yield Langmuir waves can be quite efficient.

The required large-scale magnetic field for this case, with $\omega_H = \omega_p$, is $B = 1.7 \times 10^6 \text{ G}$. Such a field is large, and does exceed the equipartition strength for the temperatures and densities we have used as can be seen from consideration of the plasma $\beta = 8\pi nk_B T/B^2 = 3.0 \times 10^{-10} T$. Nonetheless, our lack of direct information on the physical parameters in the γ -ray emitting region certainly seems to allow for such high fields as well as the high densities we posit. The required B field can be reduced somewhat if $j \geq 2$ is assumed, but then the close agreement between $L(\omega_0)$ and $L(\omega_1)$ no longer holds. For example, with $j=2$, maximum useful cyclotron

power comes from $\omega_0 = 5\omega_H$ and $\omega_1 = 3\omega_H$, and if the same volumes radiate both harmonics, $L(5\omega_H)/L(3\omega_H) \approx 0.07$. Thus the $j=1$ case is most efficient, and using it and other values from Section 4.2.2 we find at r_i (10^{15} cm) that $B_i = 1.5 \times 10^3$ G, while at r_o (10^{16} cm), $B_o = 1.1 \times 10^2$ G. Thus, in this picture, the magnetic field and density vary in tandem so that the resonance condition, $\omega_p = s\omega_H$, is satisfied throughout the volumes emitting gamma through to optical radiation. Since coherent processes can yield large luminosities from smaller volumes, they also make it easier to explain the observed rapid luminosity variability.

A condition necessary for the exponential growth of these Langmuir waves through RFS arises when one considers the balance between the forward reaction and the back reaction (absorption of a plasmon to increase the cyclotron harmonic). The rate of production of Langmuir photons is proportional to: $N_u(1 + N_l)(1 + N_p) - N_l N_p(1 + N_u) \approx N_p(N_u - N_l) + N_u N_l$. Here N_u , N_l and N_p are the occupation number of particles in the upper and lower cyclotron harmonics, and in the Langmuir wave, respectively. For rapid growth either $N_u > N_l$, or the Langmuir wave occupation number is small, which we do not want. The former condition is expressed as $E_u^2/\omega_u > E_l^2/\omega_l$, or

$$L_{s+1}/R_{s+1}^2 > (L_s/R_s^2)(s+1)/s, \quad (\text{A5})$$

where we have assumed $j=1$ and use the notation $L_s = L(\omega = s\omega_H)$, and take R as the size of the region of emission of that harmonic. Substituting for the L s we obtain from (A2)

$$\begin{aligned} & (n_{s+1}/n_s)(R_s^2/R_{s+1}^2)(N_{s+1}/N_s)(\beta_{s+1}^{2s+2}/\beta_s^{2s}) \\ &= (n_{s+1}/n_s)^2(R_{s+1}/R_s)(\beta_{s+1}^{2s+2}/\beta_s^{2s}) \\ &> [(s+1)/s][(s+1)/(s+2)](1-s^{-2})^{(2s-1)/2} \\ &\quad \times [1-(s+1)^{-2}]^{-(2s+1)/2} s^{2s+1}(s+1)^{-(2s+3)} \\ &\quad \times (2s+3)!/(2s+1)! \end{aligned} \quad (\text{A6})$$

For the conditions that concern us, $n_{s+1}/n_s \approx 1$; and, as we have discussed above, for $j=1$, $R_{s+1}/R_s \approx 1$ as well. Then we can evaluate (A6) for various values of s : for $s=2$, (A6) becomes $\beta_3^6/\beta_2^4 > 0.6028$, and if $\beta_2 = 0.3$, to use the same value as above, $\beta_3 > 0.412$; for $s=3$, (A6) yields $\beta_4^8/\beta_3^6 > 0.598$, and if $\beta_3 = 0.3$, $\beta_4 > 0.380$; $s=4$, if $\beta_4 = 0.3$, $\beta_5 > 0.362$; for $s=5$, if $\beta_5 = 0.3$, $\beta_6 > 0.350$. Thus we see that for higher harmonics the difference between the β 's decreases, and if higher harmonics are used they help in two ways: (1) one can have $N_u > N_l$ for a small difference in β_u and β_l , and (2) the original condition $\omega_0, \omega_1 \gg \omega_p$ (which here is also ω_H), is clearly satisfied and ensures that the back effect of the plasma waves on the cyclotron waves is small. Thus we have electrons with slightly higher values of β yielding the upper harmonic and those with slightly lower values yielding the lower harmonic.

Once $N_u > N_l$, the exponential growth of the Langmuir waves is guaranteed. A situation with $N_u > N_l$ can be achieved through a milder form of non-thermal distribution than that needed for masering, such as a two-temperature plasma. Such two-temperature situations are common in plasmas, e.g. in a magnetic field $T_\perp \neq T_\parallel$ because of anisotropic transport coefficients.

The use of cyclotron photons for the seed fields thus helps ensure monochromaticity and also allows for an efficient generation of continuum emission via Raman scattering. Another positive point is that the initial electron population is basically thermal and subrelativistic in this case. In order to create the desired beam we have the beating of cyclotron waves producing the Langmuir waves which then accelerate a small fraction of the electrons to large γ factors. Despite the large ambient fields, there is no problem in having electron beams develop; in fact the perpendicular motion will be radiated away as cyclotron and synchrotron photons so that the residual dominant motion is along the ambient field (e.g. Blandford & Payne 1982).