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A new Stark effect diagnostic was used to measure the E-field distribution of Langmuir waves in beam-plasma turbulence. Abrupt beam cut off causes the distribution's amplitude to decay exponentially in a microsecond, in disagreement with recent power law scalings deduced from cascade theory.

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FINAL REPORT

FOR

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ANOMALOUS DECAY OF LANGMUIR TURBULENCE

AFOSR-87-0217

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ABSTRACT

Using a Stark effect diagnostic, we measured the electric field distribution of Langmuir waves, P(E), in beam-plasma turbulence. When the destabilizing beam abruptly cuts off, the form of $P(E) \propto \exp(-E^2)$ discovered earlier persists, with amplitude decaying exponentially in a microsec. Strong fields last much longer than other time scales in strong turbulence theory. Exponential decay disagrees with recent power law scalings deduced from cascade theory. To explain this we suggest that Langmuir energy persists at long wavelengths, slowly coalescing around nucleation density wells left by previous, "burnt-out" solitons.

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Strong beam-plasma instability drives high amplitude electric fields, self-compacting density wells, and powerful microwave emission.[1-10] Descriptions of these complex phenomena envision the linearly unstable waves as the energy source, with nonlinear processes moving waves to shorter wavelengths, forming compressed soliton-like zones. These collapse to still shorter scales until they are damped by electron heating.

While this picture may yield a general scheme suitable for a wide range of differing systems, the movement of energy in wavelength and development of high electric field strengths E is cloudy. The standard picture envisions rapid, dissipation-free energy cascade in wave number k, until damping at $k\lambda_D \sim 0.1$ disperses the energy into heat and electromagnetic radiation, with λ_D the Debye length.

The Zakharov equations which govern this realm apply to fields such that $W = \langle E^2 \rangle / 4\pi nkT \ll 1$. Yet cascade often produces "cavitons" which self-confine themselves by digging density wells, reaching $W \geq 1$. Much theory and simulation explores the $W \ll 1$ regime,[3-5] but the highest electric fields and strongest electromagnetic emission occur near the end of the cascade,[11,12] $W \sim 1$. There, the power law of the probability distribution in E, $P(E) \propto E^{-b}$, with b a constant, becomes a Gaussian,[11] $P(E) \propto \exp(-E^2)$. Though power law cascade dominates k-space, the Gaussian covers most of E-space (Fig. 1).

The movement of energy is still not well studied experimentally. Though the Langmuir wave collapse paradigm governs part of the E-space picture, we do not generally know how natural beam-plasma systems transport wave energy in time. The terminating Gaussian distribution may arise from random effects at the beginning of caviton collapse,[13] or from caviton interactions. We now use this Gaussian signature distribution of the dissipative regime to study energetics in turbulence, once the driving beam is turned off. This reveals the relaxation time scales of strong wave fields without the complication of fresh energy input.

We employ the same beam-plasma system as in earlier work. We fire a 1 - μs, 15-kA, 700-keV relativistic electron beam from a vacuum diode, into a 1.5-m drift tube filled with partially ionized helium. The beam ionizes ~ 10% of the gas, producing a plasma density n_p in the range 10^{13} cm^{-3}. A 1-kG magnetic field provided beam guidance. We measured microwave emission and studied the Langmuir field strengths using a Stark effect diagnostic with optical fluorescence. The setup is exactly that of Ref. (11).

With a "crowbar" technique we can cut the 1 μs beam off in less than 50 ns at any desired time. The crowbar structure contains a trigger disk isolated from surrounding metallic parts. An isolation gap can withstand the normal trigger disk charge voltage, but arcs over when a trigger signal is applied. The trigger circuit provides 80 kV to the isolation gap, producing plasma at the cathode. Anode-cathode breakdown cuts off the 700 kV across the switch. (Fig. 2)

Optical emission from the beam-plasma interaction region travelled through a lens system into a 1 mm silica fiber, leading to a 0.75 meter monochromator. The output slit of the monochromator fitted into a fiber optics matrix, converting it into an eight channel polychromater. Eight photomultipliers transformed the optical signal into electric signals recorded by eight oscilloscopes.

As in earlier work,[11-17] we detect Stark shifts in neutral He within the beam volume. From optical fluorescence we measured the nonzero shift from rapidly oscillating Langmuir fields. We used a quadratic Stark system in singlet He, the 3^1P_1 and 3^1D_2 levels. We checked noise levels by looking at channels away from the He line with the beam on, and also by looking at the He line without the beam. Both cases gave an average of one count per ten shots, far below the photons per channel at the line shift. Null tests with the optical fiber from the blocked plasma gave the same noise level.

We counted photon numbers N(E,t) after beam cutoff (Fig. 2). The $\exp(-E^2)$ distribution discovered in Ref. (11) holds for long times after the driving beam is gone. Fig. 3 shows data collected over a μs. after beam cutoff, with mean deviations. Data from the first 500 ns. had the same distribution as in the second 500 ns. The exponential line is the best fit, though we could fit a power law, E^{-2}, to all but the highest data point in E. Figure 4 shows that the sum over electric fields, N(t), decays exponentially with time constant $\tau = 1$ μs. This is true of both low-E and high-E ends of the distribution, i.e., the mean field E does not change with time. Errors are larger at high E because of declining photon numbers.

The behavior $P(E) \propto \exp(-E^2) \exp(-t/\tau)$ suggests a selfsimilar statistical distribution decaying as energy moves to high electric fields and dissipates into plasma waves. Simple wave damping by collisions takes 10 μs (the "Coulomb" line in Fig. 4). Electrostatic wave convection from the chamber takes even longer.[16] Radiation losses from dipole emission by cavitons should scale as E^2, rapidly depleting the high field distribution and affecting the form of P(E); this is not observed.

The beam current induces return currents which persist after beam cutoff. This can drive electrostatic instabilities, since the drift velocity $u \cong n_b\, c/n_p$ is several times the electron thermal speed, allowing ion acoustic waves to grow as long as $T_e > T_i$. However, we found that the return current decayed in 10 μs., much slower than the Langmuir turbulence. This rules out return-current-driven microinstabilities as the cause of the persistent high electric field levels. (Some low-E turbulent fields do remain, causing the 10 μs. decay. This time implies an anomalous collision frequency a hundred times the Coulomb frequency. Attributing this to ion acoustic (IA) waves[18] implies $W_{IA} \sim 10^{-2}$, which is below our observable range.)

Nonlinear loss mechanisms seem more plausible, but the long microsecond decay is puzzling. Caviton collapse occurs very quickly, typically a few ns. for our experiment.[11] The simple rate equation model developed earlier for this experiment,[12] using modulational instability to transfer mode energy from the driver electrostatic instability, suggests time scales of at most 100 ns ("turbulent cascade" in Fig. 4). Turbulent processes seem too fast and Coulomb rates are too slow to explain the microsecond decay. We must then look for a reservoir of mode energy which slowly transfers its energy to the dissipating region of k.

We measured microwave emission at the plasma frequency during and after the beam pulse. Strikingly, emission dies with beam cutoff, i.e., within 50 ns. Our earlier rate equation model[12] implied that energy should cascade within 100 ns. to the "burnout" region of dissipation, so that microwaves should indeed vanish quickly. However, that model predicts that the electrostatic waves should also cascade and damp, and they do not. In view of the persistent electrostatic turbulence, we expected radiation after the beam pulse, because solitons can radiate.

To explain this puzzle we explored soliton emission during collapse. Though this is the common explanation of Type III solar bursts,[1,9,19] we found that emission by dipolar fields in solitons does not yield enough power to explain our observations. Also, we would be hard pressed to explain why the radiation at ω_p ceased, since the Langmuir turbulence causing soliton formation and collapse persists. (We could not reliably resolve the $2\omega_p$ emission expected from solitons, to see if it, too, lasted ~ μs.)

We then turned to emission by the beam itself, assisted by the strong Langmuir turbulence. This neatly explains the dying microwave emission, since beam electrons are crucial players. Relativistic electrons experience impulsive collisions with the soliton fields, launching electromagnetic wave packets of broad spectrum from the plasma frequency ω_p up to ~ $\gamma^2\omega_p$, with γ the

Lorentz factor. Most of the power lies in the plasma frequency line, as we observe. (At higher beam densities the power peak shifts to the $\gamma^2\omega_p$ region, as reported earlier.[13]) Essentially the spectrum is like that of collisional bremsstrahlung, with the soliton size L playing the role of the minimum impact parameter.[14] Emission near ω_p can escape because the plasma density declines radially. A straightforward calculation[14] yields the emitted beam power at ω_p,

$$P = \left(\frac{E^2}{8\pi} c\sigma_T\right) \left(\frac{4}{3} \pi L^3 n_b\right) \left(f n_b V\right) \frac{27\pi}{4} \qquad (1)$$

The first term is the classical Thompson scattering rate. The second is the number of beam electrons in the soliton volume, arising from the coherence of electrons passing within L of the center. A total of $n_b V$ beam electrons radiate, with V the beam volume, and solitons occupy a fraction f of this volume. We earlier determined this fraction to be between 0.05 and 0.2, using the intensity of optical Stark shifts.[15-17] We assume solitons randomly oriented in space, since we found no polarization in the microwave emission. Equation (1) yields

$$P = 30 \left(\frac{n_b}{10^{11}\ \mathrm{cm}^{-3}}\right)^2 \left(\frac{T}{100\ \mathrm{eV}}\right)^{5/2} \left(\frac{L}{\lambda_D}\right)^3 fW \quad \text{Watt} \qquad (2)$$

which is compatible with our observed power of 10 to 100 kW while the beam is on, if f is 0.1, T is 50 eV (as heating rate calculations imply[12]), $L \sim 20\ \lambda_D$ and W is of order unity. This lends support to our Stark shift measures, which find $W > 1$ for most of the distribution of Fig. (2).

This explains the quick drop in microwave emission, but why do the turbulent fields last so long? A clue emerges from re-examining our assumptions.

An earlier model for our system[12] focused on energetics, including the basic beam-plasma instability, Ohmic heating from both classical and anomalous resistivity, and electromagnetic emission. It used the oscillating-two-stream instability to transfer energy from the long-wavelength driver modes ($\lambda_o \sim$ cm). While the model gave a reasonable semiquantitative picture of the overall time evolution of turbulent energy, it also predicted a wave spectrum which rose exponentially in k until the damping regime set in at $k\lambda_D \sim 0.1$. This odd outcome may arise from the one-dimensional model, since phase space volume effects yield steeper descending spectra in higher dimensions. (Beam effects probably do not produce large anisotropics in the plasma environment of our experiment, though, because the unstable wavelengths greatly exceed the critical wavelength ($\sim \lambda_D (9M/4m)^{1/2}$, with M/m the ion-electron mass ratio) above which backscatter cascade via three-wave decay is impossible.)

In any case, the predicted increasing spectrum disagreed with all previous theory and simulations. Further, it failed to predict the persistence of turbulence after beam cutoff.

Recent work shows that the oscillating two stream instability and modulational instability require quite uniform plasma, or else they are swamped by the local forming of coherent wave packets.[8,20] Density variations arise from the burnout of previous collapsed wave packets. Nucleation around these burnout sites is faster than the oscillating two stream, et. al., and yields a decreasing power-law probability P(E) up to the characteristic wavelength where dissipation sets in. Beyond that point P(E) becomes, as we observe, $\exp(-E^2)$. (Our efforts to measure the power law portion of the spectrum at lower E have so far failed due to difficulty resolving small Stark-effect splittings.)

Nucleation proceeds as energy flows from a sea of background states into localized ones, even without a driver. Robinson and Newman[20] recently proposed a model with energy not entering directly from the driver fields. Instead, energy flows from background turbulence into wave packets which form in the shallow density wells left from burnout. This suggests a bank of wave energy in the background states, approximately plane-wave modes, which can take long times to focus on a density well -- longer than later collapse and dissipation. Once the driver (beam) stops, this bank bleeds into the localized eigenstates at a rate difficult to predict. This may explain our μs. decay.

Changes in the density wells couple the localized modes to the approximately plane wave background states. (Absent such changes, the states are orthogonal and noninteracting.) Energy will flow as long as the density wells do not relax so fast that the bound eigenstates vanish before the local field strengths reach the threshold for collapse.

Sound waves can disperse wells by transporting energy, relaxing ion density gradients, in a time $\sim L^*/C_s$, with L^* the size of wells before they reach threshold, and C_s the sound speed. Robinson & Newman[20] argue that the threshold size of wells, L_t, must exceed $(M/m)^{1/2}\ \lambda_D$. For our experiments in helium, $L_t/C_s \gtrsim 0.1\ \mu s$. This argues for a longer decay time than is available in any cascade scheme, though it is difficult to fix an upper bound. The largest possible length for a density well is the scale of the linear instability, $\lambda_0 = 2\pi\ v_b/\omega_p \sim$ cm. Sound waves could relax such a scale in $\sim 5\ \mu s$.

For many purposes, snapshots of the underlying turbulence are not germane in describing naturally arising turbulence. Instead, statistical properties, spectra and scalings, for averages over the distribution of many solitons in a statistical steady state, are more interesting and useful. To this end the exponential time decay of N(t) is significant. Most turbulence models yield power law decays in time as well as in k.

As the mean dimensionless energy $\langle W \rangle$ rises, the $\exp(-E^2)$ portion of P(E) contributes a larger fraction to $\langle W \rangle$ than that from the power law at lower E. For the large $\langle W \rangle \gtrsim 1$ we observe, the time evolution of $\langle W \rangle$ may be dominated by the behavior of the high-E tail. In the scaling model of Robinson & Newman,[20] $\langle W \rangle$ scales as $t^{-2/D}$ with D the dimensionality. This differs considerably from our exponential.

Earlier work deduced a power law in P(E) for collapse.[1,4,5,19] The Robinson and Newman model modifies this to produce an exponential tail, essentially by assuming random initial conditions in the collapse. This still yields a power law in the time decay after the driver shuts off. Our observation of exponential decay implies an important missing element in existing models. Energy transfer may depend significantly on the first, slow step in the overall process, before caviton formation. Probably a detailed study of energy transfer between the background turbulence and localized wells can give an exponential, since any random process of localization will suffice, though the long μsec decay time is still a striking problem.

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FIGURE CAPTIONS

Fig. 1. Schematic of strong turbulence dimensionless energy W(k) in k-space. P(E) is the corresponding probability of observing electric field intensity, E. Background turbulence (low k, shaded) contributes to low E. After nucleation around shallow density wells, energy flows to higher k and higher, localized E, during cascade. Dissipation (second shaded region) changes the power law behavior of W(k) and P(E).

Fig. 2. Typical observations during a beam pulse. Diode voltage (a) and beam current (b) cut off abruptly, as does microwave emission in two representative channels (c,d). Optical emission in Stark-shifted lines (e-h) shows persistent photon counts for a μs. after beam cutoff.

Fig. 3. The number of photons observed at differing Langmuir field levels, E. These fields oscillate at the plasma frequency and obey an exp $(-E^2)$ distribution for all times after beam cutoff. This data is for the first 500 ns. after cutoff.

Fig. 4. Decay of the total integrated intensity of photons, N(t), after integration over E. Beam cutoff is at t=0. Coulomb collisions yield a ~ 10μs decay time. Cascade in k-space should occur in ~ 0.1 μs.

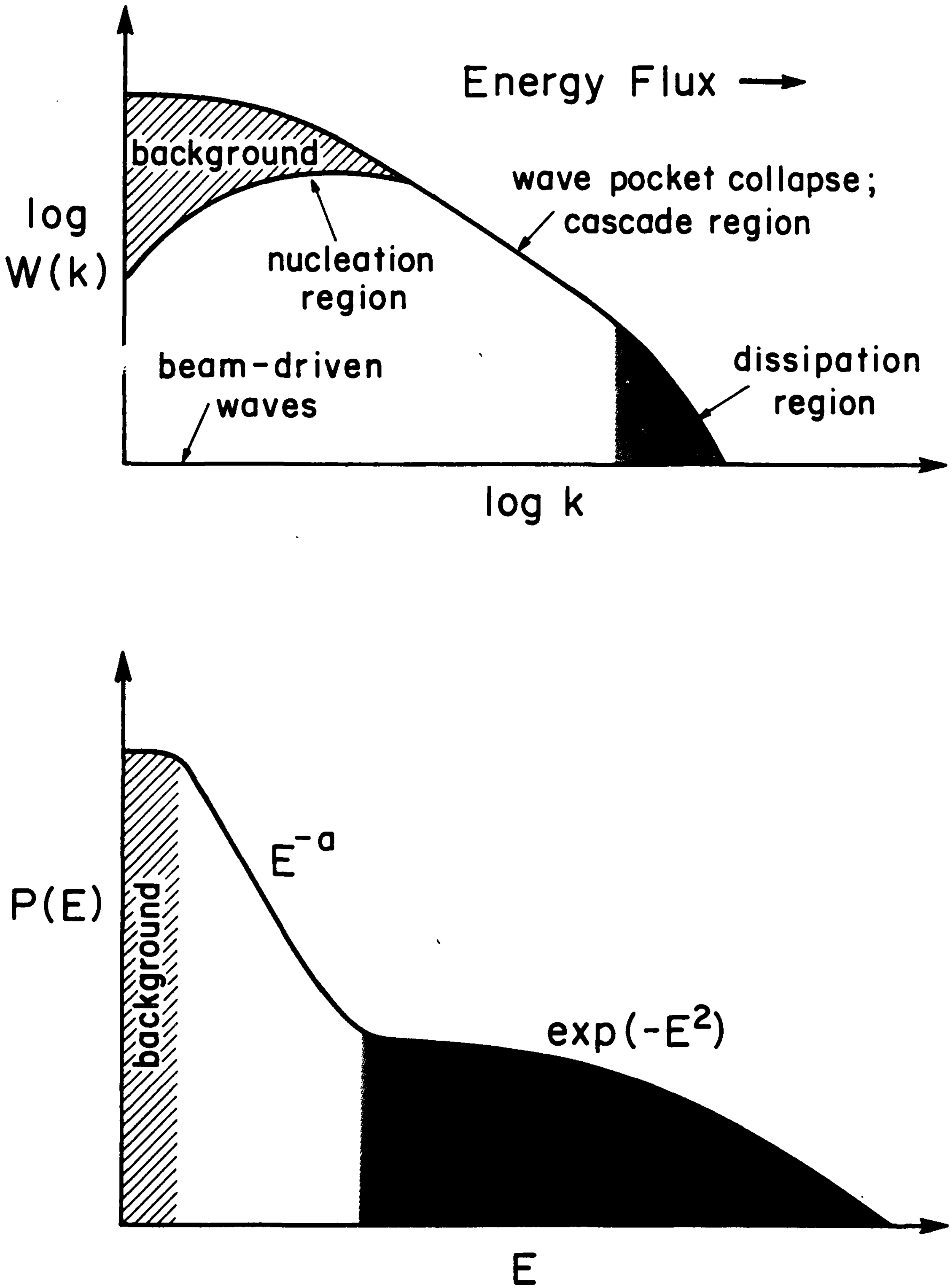

Energy Flux →
log
W(k)
background
nucleation
region
wave pocket collapse;
cascade region
beam-driven
waves
dissipation
region
log k
P(E)
background
E^-a
exp(-E^2)
E

Figure 1

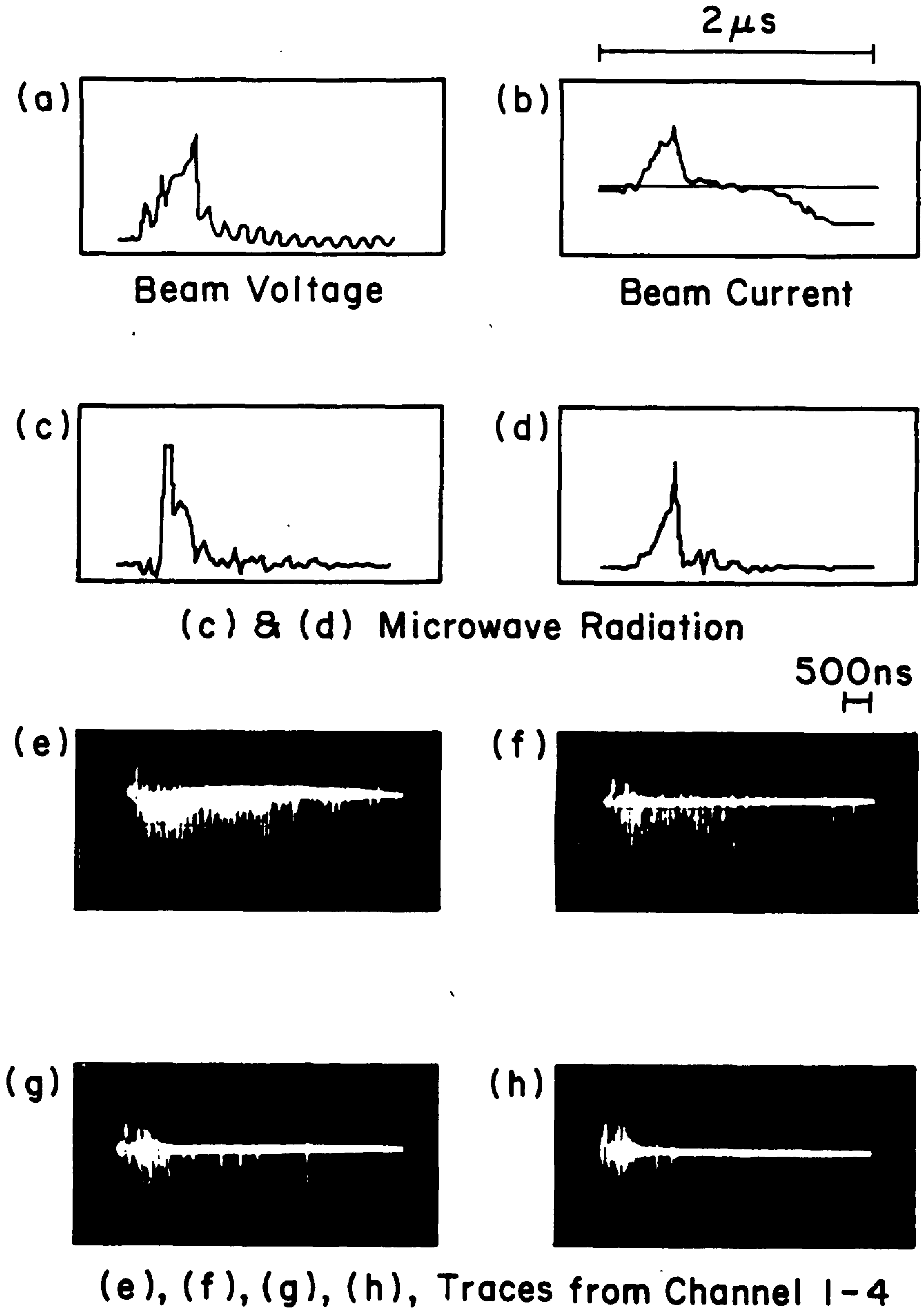

2μs
(a)
Beam Voltage
(b)
Beam Current
(c)
(d)
(c) & (d) Microwave Radiation
500ns
(e)
(f)
(g)
(h)
(e), (f), (g), (h), Traces from Channel 1-4

Figure 2

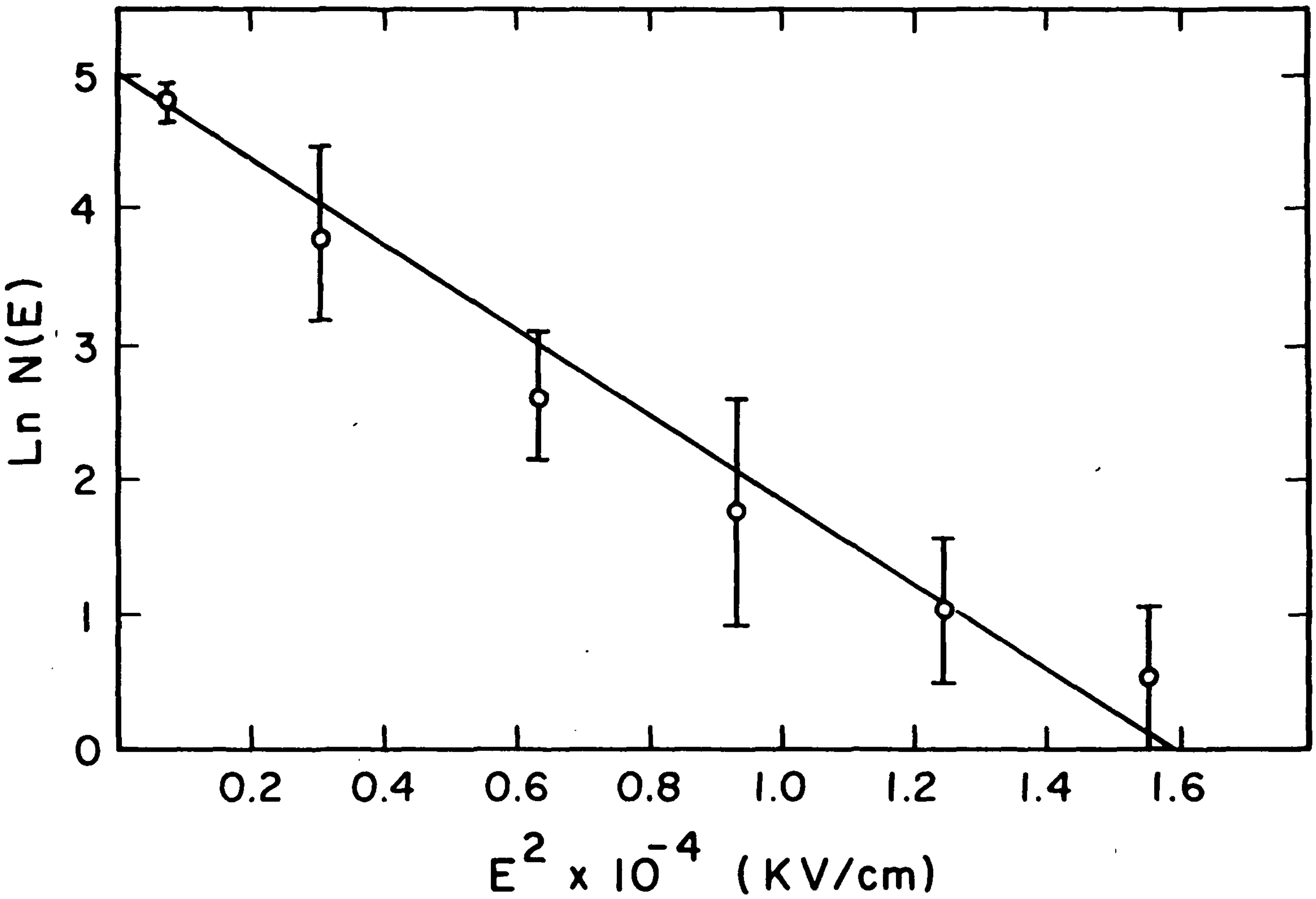
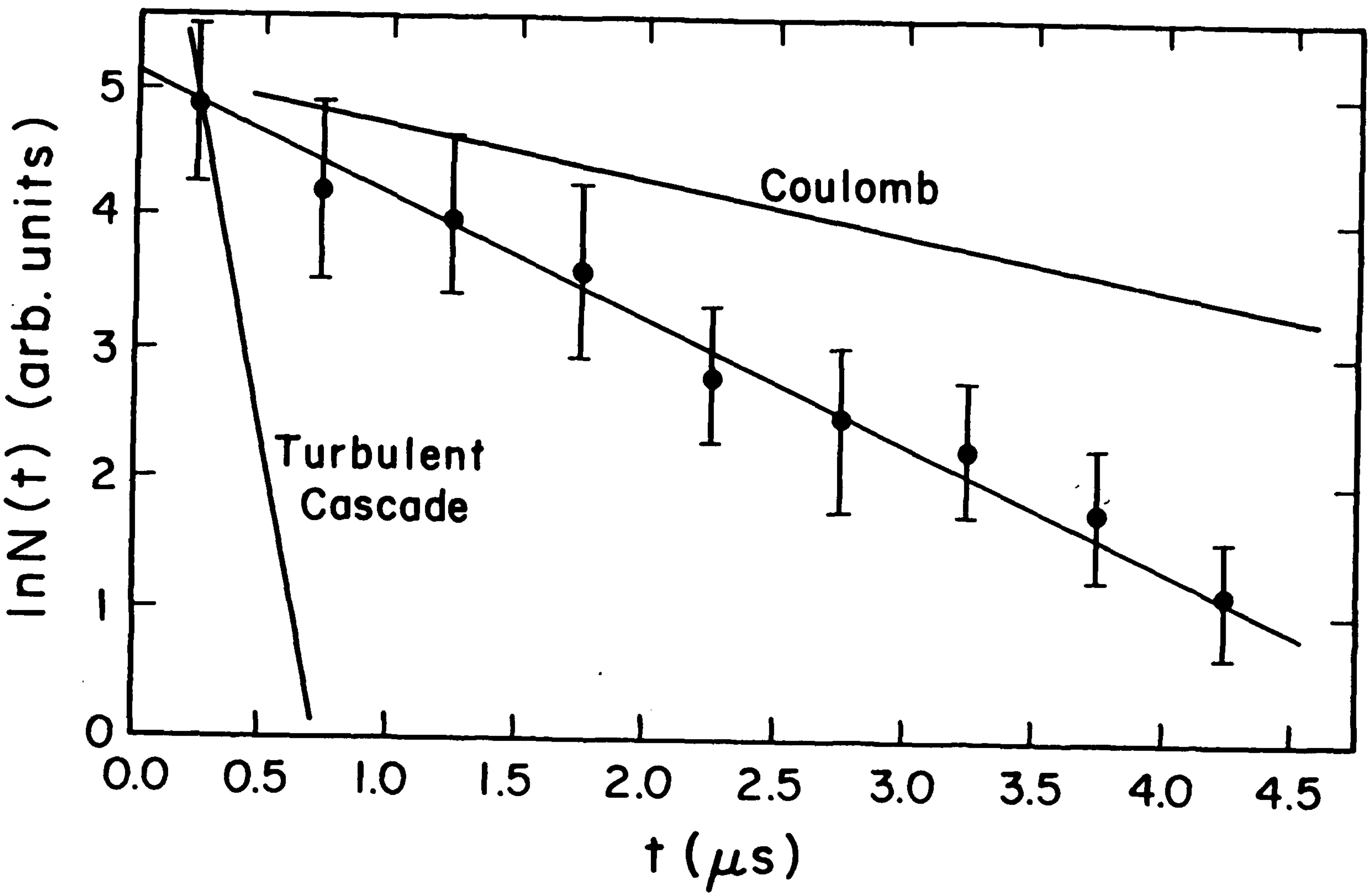


Figure 3



lnN(t) (arb. units)
5
4
3
2
1
0
0.0
0.5
1.0
1.5
2.0
2.5
3.0
3.5
4.0
4.5
t (μs)
Coulomb
Turbulent
Cascade

Figure 4