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ON THE POSSIBILITY OF COHERENTLY STIMULATED RECOMBINATION AND COSMOLOGICAL STRUCTURE GENERATION *

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Abstract

Possible instabilities during cosmological recombination may produce an epoch of non-linear density growth and fractal-like structural patterns out to the horizon scale at that epoch ($\sim$ 200 Mpc today). Such effects could explain observed large-scale structure patterns and the formation of objects at high z while keeping microwave background anisotropies at minimal levels. With this motivation, we examine the consequences of the change in effective radiative recombination reaction rate coefficients produced by intense stimulated emission. The proton-electron recombination is considered as a natural laser, leading to the formation of spatially non-uniform distributions of neutral matter earlier than the recombination epoch. A model for such fractal patterns is presented. We also discuss possible microwave background implications of such a transition and note a potentially observable spectral signature at $\lambda \sim 0.18$ mm as well as a weak line near the peak in the microwave background.

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INTRODUCTION

A forefront problem in cosmology today is how to generate the observed large-scale structure and allow the existence of objects at high redshift, z, without requiring unobservable fluctuations in the cosmic microwave background radiation (CBR). One avenue that has been proposed to surmount these problems has been a late-time vacuum phase transition occurring just after recombination (Hill, Schramm, & Fry 1989). We note here that non-linear growth of instabilities at recombination itself can provide many of the benefits of a late-time vacuum phase transition without having to postulate a new fundamental interaction occurring at an apparantly arbitrary energy scale.

Although earlier work (Hogan 1989) had been pessimistic about the possibility of instabilities at recombination, several new developments now offer encouragement to this possibility. Of greatest significance has been the recent actual experiments of Schramm et al. (1991) and Yousif et al. (1991) showing that laser-stimulated recombination does indeed occur with cross-section gains of several thousand. Furthermore, Hogan's initial pessimism was based on the assumption of spherical geometry for the recombination sites, whereas we argue here that the geometry is auto-catalytic in nature and will produce a fractal-like structure somewhat analogous to the diffusion-limited aggregation models of Peters et al. (1979), possibly with the "shadowing" effects of Witten & Sander (1981). Also, Hogan (1991) has recently noted that the instability in the trapping of Ly-α photons can at least yield fluctuations of significance to star formation size seeds. We argue here that the new experimental, laser-stimulated recombination results may even suggest a much grander scale of instability that could account for galaxies and large-scale structure without seriously disturbing the microwave radiation beyond its observed limits.

To appreciate the possible significance of an instability of recombination, let us note that a non-linear instability at recombination could yield a fractal pattern of neutral gas on scales comparable to the horizon at the time of recombination ($R_{comoving} \sim 200$ Mpc). The neutral gas regions can gravitationally contract into objects while the ionized gas remains uncollapsed. Hogan's (1991) cosmic Doppler instability could further enhance the gravitational growth. It is even possible that the rapid collapse of neutral regions and the subsequent formation of radiating objects could prevent the remaining ionized regions from ever recombining. This could explain the absence of a Gunn-Peterson trough in the spectra of high redshift quasars (Schneider, Schmidt, & Gunn, 1991).

The fractal patterns could yield the observed large-scale patterns out to scales $\sim$ 200 Mpc and might be consistent with the Szalay-Schramm (1985) interpretation of the cluster-cluster correlation of Bahcall & Soniera (1984) and Efstathiou et al. (1991).

While the above possibilities are somewhat speculative, they do show some potential roles that a recombination instability could play.

In this paper we explicitly suggest the feasibility of the occurrence of a strong non-linear coupling between radiation and the proton-electron recombination kinetics. The non-linearity produced by this coupling must be spatially non-uniform and therefore leads to a spatially non-uniform distribution of product.

We will first describe the appropriate atomic physics; then we will present a simple model for the possible patterns produced; and finally we will discuss possible microwave background implications.

COHERENTLY STIMULATED RECOMBINATION

In simplest form recombination is viewed as the process

$$p + e = H + h\nu. \tag{1}$$

The evolution of the $\frac{p}{H}$ ratio is conventionally viewed as essentially determined thermodynamically or by equivalent equilibrium statistical arguments, once the baryon/photon number is set. As has been discussed frequently (Krolik 1989), the validity of this treatment depends upon the kinetics of the processes involved. Each epoch is unique since evolution is not cyclical. Thus, the validity of the evolutionary arguments depend critically upon the accuracy and most importantly the completeness of the kinetic considerations. Here we suggest the possibility of new kinetics that would occur in the presence of intense coherent radiation.

We argue that in a region exposed to a fluence beyond a critical value, the proton-electron recombination kinetics will be essentially radiatively stimulated. The rate of this process has a modest temperature dependence and is much faster than the rate of spontaneous recombination. The reverse process, ionization of the neutral hydrogen, will still occur, as in all recombination models, by thermal cosmic background radiation. The rate for this process is a strong function of temperature. Under these circumstances the $\frac{p}{H}$ ratio will not be that determined by the temperature of the cosmic background radiation in all spatial regions, but will depend upon the structure of the propagation of the pulses of coherent radiation.

As is well established, the rate of transition between bound states of hydrogen may be determined by the density of coherent radiation, either by one photon or multiphoton resonant processes. The reaction $p + e = H$ depends upon the transition from a free to a bound state. It has recently been shown that the reaction rate may be enhanced by orders of magnitude by coherent radiation of modest intensity (Schramm et al. 1991; Yousif et al. 1991). This combination of processes, together with rapid transitions between high n states induced by the cosmic background radiation, provides a rapid route for the condensation $p + e = H(1s)$. For purposes of clarity, we illustrate the proposed kinetic path with a somewhat specific example, acknowledging that there may be alternate routes which are superior in efficiency. The parameters used are

$$\frac{n_\gamma}{n_b} = 3.3 \times 10^9$$

$$T_0 = 2.73$$

$$n_\gamma = 411$$

$$\Omega_b = 0.06$$

$$n_0 = 1.25 \times 10^{-7}$$

We make estimates at Z=1500.

Figure 1 illustrates schematically the process. Electrons from the thermal distribution undergo stimulated transitions to $n \geq 6$. The coherent radiation field is a pulse of order $10^9 \mathrm{Jcm}^{-2}$ fluence, with frequency $\frac{1}{16} R_H$. These levels are very appreciably Stark broadened (Autler & Townes 1955) so that virtually all of the free electrons states are connected to the bound states by a one photon process. The high n levels are resonant with the n=4 level which is Stark broadening to approximately 200cm^{-1}. A consequence of the strong electric field is to destroy angular momentum as a good quantum number. Thus, the nominal selection rule producing only high n states of high angular momentum, i.e. $l \approx n$ in the initial electron capture, is followed by rapid mixing, allowing the transition to n=4. This fast, one-photon process depletes the 200cm^{-1} wide resonant high n distribution which is refilled rapidly by transitions induced by the continuum cosmic background radiation (Spencer et al. 1982; Kleppner 1991).

The n=4 level has a mean natural lifetime near 10^{-7}sec. The coherent radiation field is chosen large enough, by specifying the pulse width that the three photon resonant n=4—n=2 transition occurs in 10^{-8} sec. The coherent radiation field as specified has an energy density of three orders of magnitude greater than the cosmic background radiation. As is conventional, the n=2 level decays by a two-photon process. In our model, however, this process consists of one coherent—one spontaneous rather than the normal two spontaneously emitted photons. Our estimate for this time is 10^{-4}sec as contrasted to $\frac{1}{7}$sec for the spontaneous 2s—1s transition. A virtue of this process is that the frequency of the spontaneous photon insures energy conservation in the overall process. The spontaneous emission is further required to give the unidirectional character of the process.

Examination of the lifetimes of states with respect to ionization by the cosmic background radiation shows that at redshift z=1500, all of the coherent radiative processes of a level occur more rapidly than ionization. In addition, the ground state neutral hydrogen is stable, while at z=2000, its stability is problematic at best. A choice of z=1000 is inconsequential for stimulated recombination, since the plasma is already neutral by spontaneous recombination.

The intrinsic bandwidth of the laser described above is likely to lie between the Stark broadening of the n=4 and the n=2 level. An estimate of the former gives $\frac{\Delta\nu}{\nu} = 0.03$, while the latter gives $\frac{\Delta\nu}{\nu} = 0.007$. At any time, in any specific region, the monochromaticity is likely much higher. The linewidth of the incoherent photon with frequency near $\frac{11}{16}R_H$ is somewhat greater than that of the coherent photon. It is this radiation which undoubtedly is the most observable signature of this postulated event. At z=1500 for this radiation which contains more than half the recombination energy, $x = \frac{h\nu}{kT} = 26.5$. Since the blackbody spectrum is extremely low, almost irrespective of the extent of broadening, this should be discernable in the neighborhood of $50\text{cm}^{-1} = 200\mu$. We have chosen pulse parameters somewhat casually. All of the atomic processes are fast with the exception of the frequency of electron proton collisions. The rate of this process depends upon the extent to which high n=l states are mixed with $l \leq 4$ levels, allowing transition to the n=4 level to occur rapidly. Thus, it is likely that the cross section for the initial step of electron capture sets a minimum regime of time. In our estimates we have chosen a pulse of 220 sec duration. All other steps favor high fields and short times for efficiency. We have chosen a fluence as $\frac{6}{17} \times \frac{1}{2} \times (n_o z^2 \frac{\hbar}{H})(R_H + \frac{3}{2}kT)$. The gain medium is regarded as active, with the radiation pulse having an appreciable coherence length. As a consequence, we do not consider Thomson scattering as a photon loss mechanism. We have not discussed the difficult question of the starting of the laser. The frequency chosen has one photon resonances to a bound short lived level. as well as the multiphoton resonances. A fluctuation in cosmic background radiation may therefore be amplified at low field strength. We have not examined the likely coherence length in the radiation pulse to see whether cooperative effects at the density extant may be a factor, although, in view of the broadening of the levels, we doubt that that these are of importance.

FRACTAL PATTERNS

Figure 2 shows typical results from a numerical model of structure formation we might expect this recombination instability to produce. The numerical calculation presented here is a discrete model of the random growth of the unstable medium. We adopt a regular (three-dimensional) cubic lattice, and regard each lattice site as a potential site for instability. At the initial time, we assume that one lattice site becomes unstable. Because of the autocatalytic nature of the instability, the probability of instability of neighboring sites is now higher than that for a random "field" lattice site; hence, we chose at random one of the lattice site neighbors of the seed to become unstable. This process is continued, with the subsequent lattice sites always destabilized at random, but chosen so that they lie next to already destabilized lattice sites. This model for the spatial growth of the instability strongly resembles a model for diffusion-limited aggregation (DLA) propounded by Peters et al. (1979) in which particles are added to the growing cluster at random, but always adjacent to already occupied sites.

We note that in the limit of large cluster size, this model has the property that density correlations in the recombined phase are independent of distance.

We further note that when the "cluster" becomes sufficiently large, there may well be effects akin to "shadowing" found in DLA, though the physical mechanism is very different. In DLAs, shadowing occurs because parts of a cluster begin to inhibit or block the aggregation on "interior" sites. In our case, what can happen is that unstable material is "used up," so that lattice sites adjacent to "voids" can no longer grow. In other words, there may not be enough ionized matter to provide the "gain" necessary for the instability. In the case of significant shadowing for DLA, a more appropriate model is that of Witten & Sander (1981), in which dendridic growth at the tips of the growing cluster dominates; this gives less compact, more "wispy," structures for DLA, and it is likely (though unproven) that similar geometry will prevail in the present case. In any case, at the small cluster sizes that have been modeled so far, "shadowing" effects are irrelevant. However, a major effect ignored here is the likely anisotropy of the radiation field associated with the instability. This anisotropy (which leads to a preferred axis of the radiation field along the direction of maximum gain) would also favor elongated shapes for the recombined regions, and hence dendridic growth of the "cluster."

The patterns developed here should be appropriate to the spatial patterns of the neutral gas in the ionized plasma. Gravity coupled with Hogan (1991) processes will subsequently enhance the density in these neutral regions and produce objects. In a subsequent paper we will explore the hydrodynamics and evolution of these regions. However, for this paper we merely note the intriguing patterns that the instability may produce and speculate that subsequent structure formation may approximate these patterns.

POSSIBLE MICROWAVE BACKGROUND EFFECTS

The study of the possible deviation of the CBR from a perfect black-body spectrum

$$i_0(\nu) = \frac{x^3}{e^x - 1}, \; x = \frac{h\nu}{KT}, \tag{2}$$

goes primarily along two lines: One is the global distortion of the spectrum. It was shown (Sunyaev & Zeldovich 1970) that when there is energy release at $z > 4 \times 10^4/\sqrt{\Omega_b}$, the spectrum will exhibit μ-distortion: the background photon acquires a non-zero chemical potential μ, so the flux spectrum is

$$i(\nu, \mu) = \frac{x^3}{e^{\mu + x} - 1}. \tag{3}$$

When the energy release occurs at $z < 4 \times 10^4/\sqrt{\Omega_b}$, the spectrum can exhibit y-distortion (Zeldovich & Sunyaev 1969) due to the compton scattering of the microwave background with the hot electons. However, for our process, compton scattering is not the dominant loss mechanism and does not significantly affect the frequency. Recent COBE observations find $y < 10^{-3}$. This result has already put a strong constraint on some structure formation scenarios (Levin, Freese, & Spergel 1991) but does not constrain our mechanism directly.

A second distortion can come directly from the cosmological emission lines in the CBR spectrum. In general, during the recombination period, the excited hydrogen atom will produce line features in the relic CBR spectrum. Thus, the observation of these lines will not only provide us more detailed information about this event, but will also confirm the relic nature of the CBR. Several authors (Lyubarsky & Sunyaev 1983; Zeldovich et al. 1969; de Bernardis et al. 1990) have investigated the problem of seeing lines in the CBR and found that, for normal recombination, the intensity of line emissions is extremely low; and they are all embedded in the galactic dust emission noise—to such an extent that it will be difficult to detect them in the near future.

However, the coherently stimulated recombination scenario discussed above may enhance the possibility of success in the search for the cosmological line emission. The basic idea of the CSR is to treat the electron-proton recombination as a natural laser. The energy pump of the laser is the energy release from recombination. The working frequency of the laser is still a free parameter. By the argument of maximal energy gain, $h\nu_0$ is chosen to be $R_H/16 = 0.85eV$ or $E_{Lyman\alpha} - R_H/16 = \frac{11R_H}{16} = 9.35eV$. This corresponds to a possible spike in the CBR spectrum at wavelength $\lambda \sim 0.18$ mm or 0.20 cm. (Note that the peak wavelength of 2.75 K black-body radiation is ~ 0.19 cm.) The intensity of the spike is calculated below, and it depends on another dramatic consequence of the scenario. The occurrence of coherently stimulated photon emission strongly suggests a strong non-linear coupling between radiation and molecular kinetics. The consequence is that the chemical evolution time scale can be approximately eight orders of magnitude faster than those of the traditional homogeneous kinetic scheme. In the homogeneous case, the duration of recombination $\frac{\Delta Z_{rec}}{Z_{rec}}$ is often taken to be $0.1 \sim 0.2$, while in our case, this quantity can be as small as 10^{-9}.

A coherent photon emitted at frequency ν_0 with band width $\Delta\nu$ has an energy density spectrum which is approximately a gaussian:

$$E(\nu) = A \cdot e^{-(\frac{\nu-\nu_0}{\Delta\nu})^2}. \tag{4}$$

Total recombination per volume is $n_b E_0$. The energy release for each $e-p$ pair is $E_0 \sim 13.6eV$. If the efficiency of the energy pump is $\alpha, 0 \leq \alpha \leq 1$, then the total energy density of the coherent photon is $\alpha n_b E_0$, from

$$E_{tot} = \int E(\nu)d\nu \approx A \int_{-\infty}^{+\infty} e^{-(\frac{\nu-\nu_0}{\Delta\nu})^2} d\nu, \tag{5}$$

we have

$$A = \frac{\alpha n_b E_0}{\sqrt{\pi}\Delta\nu}. \tag{6}$$

The observed flux $I(\nu)$ is related to the energy density by $I(\nu) = (C/4\pi)E(\nu)$, so the coherent photon flux is

$$I_c(\nu)d\nu = \frac{\alpha n_b E_0 C}{4\pi^{3/2}\Delta\nu} e^{-(\frac{\nu-\nu_0}{\Delta\nu})^2} d\nu. \tag{7}$$

At the same time, the background flux is

$$I_r(\nu)d\nu = (C/4\pi)(\frac{n_\gamma}{2\zeta(3)})h\frac{(h\nu/KT)^3}{e^{h\nu/KT}-1}d\nu \tag{8}$$

where $2\zeta(3) = \int_0^\infty x^2/(e^x-1)dx \sim 2.40$. By using the same x parameter as above, the ratio of coherent photon intensity to the background photon intensity is given by:

$$\frac{I_c}{I_r} = (2\zeta(3)/\sqrt{\pi})(\frac{n_b}{n_\gamma})(\frac{\alpha E_0}{x_0})(\frac{x_0}{\Delta x})\frac{e^x-1}{x^3}e^{-(\frac{x-x_0}{\Delta x})^2}. \tag{9}$$

Our choice of parameters is: $\eta = \frac{n_b}{n_\gamma} = (3-5)\times 10^{-10}$, the efficiency coefficient $\alpha \sim 1$ at maximum gain, the recombination temperature $KT \sim 0.318eV$ ($z \sim 1300$), $x_0 = h\nu_0/KT$, $E_0/KT \approx 42.8$. Thus, the intensity ratio is

$$\frac{I_c}{I_r} = (2-3.3)\times 10^{-9} \times \frac{x_0}{\Delta x} \text{ for } h\nu_0 = R_H/16 \tag{10a}$$

$$\text{and} \frac{I_c}{I_r} = (1.7-2.8)\times 10^{-1} \times \frac{x_0}{\Delta x} \text{ for } H\nu_0 = 11R_H/16 \tag{10b}.$$

The relative intensity depends crucially on the width of the line. If there is no dramatic energy release after the recombination. the following are the possible sources that contribute to the line width.

(1) Natural width of the line.

The life time of the excited state is of order $10^{-8}s$, $\Delta\nu \times \Delta t \sim 1$, so $\Delta\nu \sim 10^{8}s^{-1}$, or $\Delta\nu/\nu_0 \sim 10^{-7}$, which is negligible.

(2) Broadening due to H-H collisions.

The mean free path of H is $l \sim 1/\sigma n_b, n_b \sim 100/\mathrm{cm}^3$ at recombination, the collision time scale $t \sim l/v \sim 10^6 s$, so collisions have almost no effect on line broadening.

(3) Doppler broadening due the thermal motion of the H atom.

The mean velocity of H atoms $\bar{v} \approx \sqrt{3KT/M_H}$, so the Doppler broadening is

$$\Delta\nu/\nu_0 \sim \bar{v}/C \approx 3.2 \times 10^{-5}. \tag{11}$$

(4) Stark broadening.

The energy of the hydrogen atom is Stark broadened by the strong electric field of the laser beam. If we choose E = 600,000 Volts/cm as a critical coherent field, then the Stark width of $n = 2$ is $\delta\nu/\nu = 0.007$ and $n = 4$ is $\delta\nu/\nu = 0.03$. The intrinsic bandwidth of the laser due to Stark broadening is likely to lie between these values, so we choose

$$\left(\frac{\delta\nu}{\nu}\right)_{Starkbroadening} = 0.01 \tag{12}$$

(5) Compton broadening due to the thermal motion of the electrons.

The Compton scattering of the photon off the moving electrons will broaden the line feature, the result of which is

$$\Delta x/x \approx \sqrt{4y} \tag{13}$$

where y-parameter is the compton parameter defined in Zeldovich & Sunyaev (1969). During recombination, the temperature of the electron is about the same as the photon temperature. $T_e \approx T_\gamma = T_0(1+z)$; $T_0 = 2.75K$, $\sigma_T = 6.65 \times 10^{-25}\mathrm{cm}^2$,

$n_e = \eta n_\gamma^0(1+z)^3$, $n_\gamma^0 = 422cm^{-3}$. At recombination, the universe is already matter dominated, so the expansion factor $\frac{dt}{dz} \approx -\frac{1}{H_0(1+z)^{5/2}}$, where H_0 is the Hubble constant. This yields $y \approx (\frac{KT_0}{m_e C^2})(\sigma_T \eta n_\gamma^0 / H_0) \int_{z-\Delta z}^{z} (1+z)^{3/2} dz$. Putting in the constants, we obtain $y \sim 3 \times 10^{-6} (\Delta z/z)_{rec}$, which yields

$$\Delta x/x \sim 4 \times 10^{-3} \sqrt{(\delta z/z)_{rec}} \tag{14}$$

which is small compared to eq. 12.

(6) The duration of the recombination.

Because the line emitted at different times will be differently redshifted, we have

$$\Delta x/x \sim \Delta z/z. \tag{15}$$

Under the conventional picture of recombination, the important contribution comes from (5), $\Delta z/z \sim 0.1$; thus, from (10a), the intensity ratio of the line to the CBR $\sim 10^{-8}$. However, in the CSR scenario, on the one hand, the reaction rate is increased by eight orders magnitude. The dominant effect is no longer (5), but the Stark broadening instead, which is one order of magnitude smaller. Putting the line width into (10), we have

$$\frac{I_c}{I_r} \approx (2.0 - 3.3) \times 10^{-7} \tag{16}$$

which is still too small to be detected in the near future. On the othere hand, CSR at $h\nu_0 = 11R_H/16$ will affect the CBR spectrum in the short wavelength region. From (10b),

$$\frac{I_c}{I_r} \approx 20, \tag{17}$$

at wavelength ~ 0.18mm..

We note here that at short wavelengths, the CBR spectrum is polluted by "dust" and other "local" sources. If future experiments such as the DIRBE experiment on COBE could subtract these local sources below the CBR level, a "spike" in the far submillimeter region would be the clear indication of coherent stimulated recombination.

The line distortion mentioned above can be a clear signature of the instability processes discussed. Because significant density contrasts probably do not develop until after decoupling, the spatial anisotrôpies generated by this mechanism will be analogous to those generated in a late-time vacuum phase transition (Schramm 1991; Hill, Schramm, & Widrow 1990).

As shown by Hill et al. (1990), these tend to yield the smallest average $\frac{\Delta T}{T}$ across the sky possible for a given size observed structure. However, as shown by Turner, Watkins, & Widrow (1991), such topological seeds can yield potentially observable spikes. A detailed calculation of the actual patterns and amplitudes for the processes discussed in this paper is being carried out now. This calculation utilizies the fractal model of Figure 2 and dynamically follows the gravitational clumping of the neutral matter enchanced by the Hogan processes. We will discuss its results in a subsequent paper and merely note here our anticipation of the similarites to late-time transitions.

CONCLUSION

We have shown that a speculative but plausible process at recombination can allow instabilities to grow non-linearly which could spontaneously yield cosmological structure. We argue that such structures might be distributed in fractal-like patterns on scales comparable to the co-moving horizon at recombination ($\sim$ 200 Mpc) and be incoherent on larger scales. Such a pattern may be a reasonable fit to large-scale structure observations. We also show that such instabilities may produce a potentially observable signature in the microwave background spectra. More detailed modeling of the spatial anisotropies of the CBR and galaxy formation modeling in this scenario will be discussed in a future paper.

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FIGURE CAPTIONS

Figure 1. Multiphoton recombination. The coherent frequency is $h\nu_0 = \frac{R_H}{16}$. The incoherent frequency ν_i is determined by $h(\nu_i + 6\nu_0) = R_H + \varepsilon$. ε is the initial energy of the electron. The Stark broadening of the $n \gtrsim 10$ levels is strongly understated for clarity.

Figure 2. A display of a discrete fractal cluster produced by the autocatalytic instability discussed in the text. The structure shown was "grown" on a cubic three-dimensional lattice, and is composed of 116 unstable (i.e., recombined) lattice sites.

Figure 3. The line features in the CBR. The width of the line is $\delta\nu/\nu \sim 0.01$; the intensity of the CBR is reduced by a factor of 2×10^{-6}.

ADDRESSES

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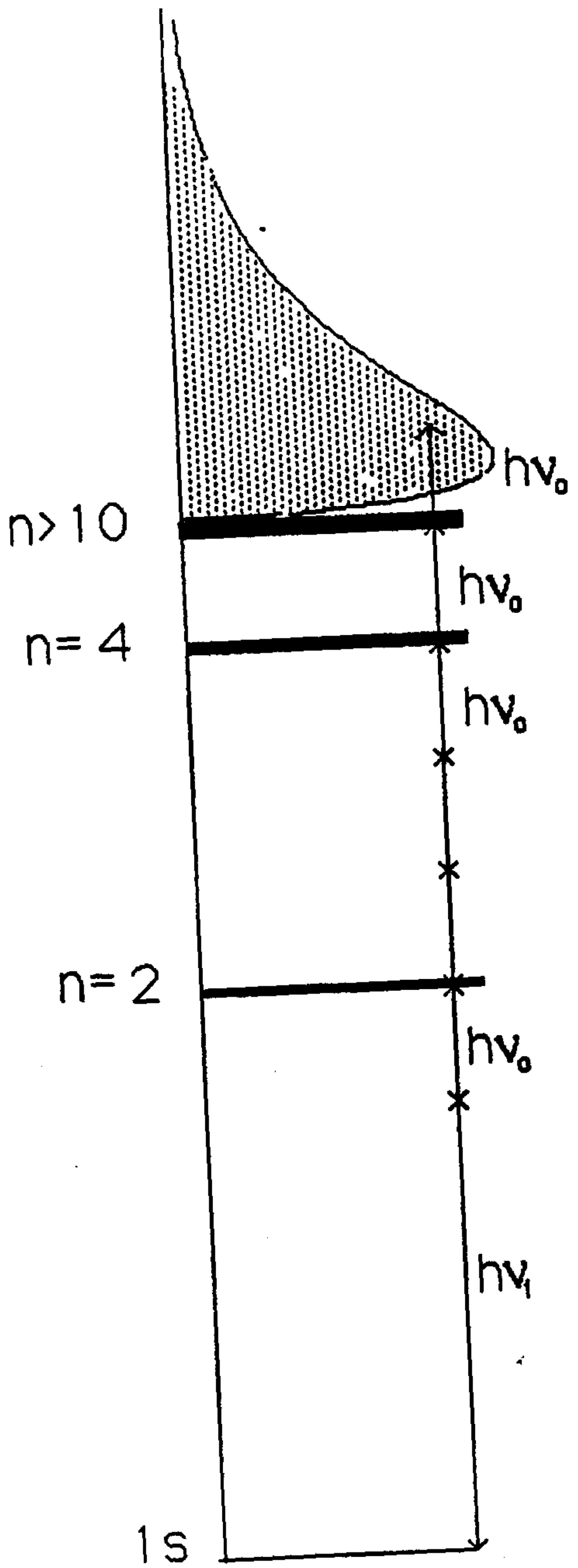

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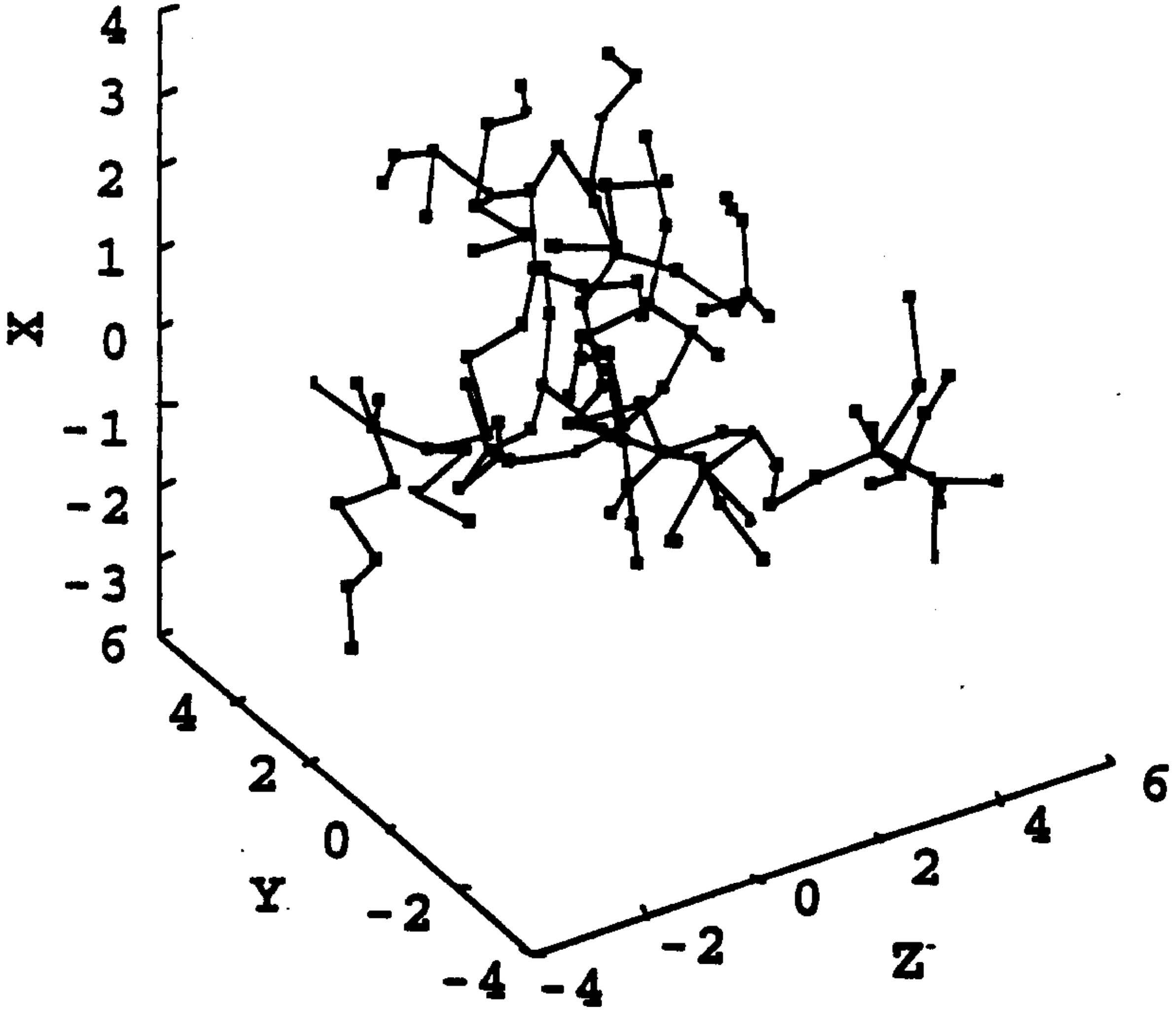

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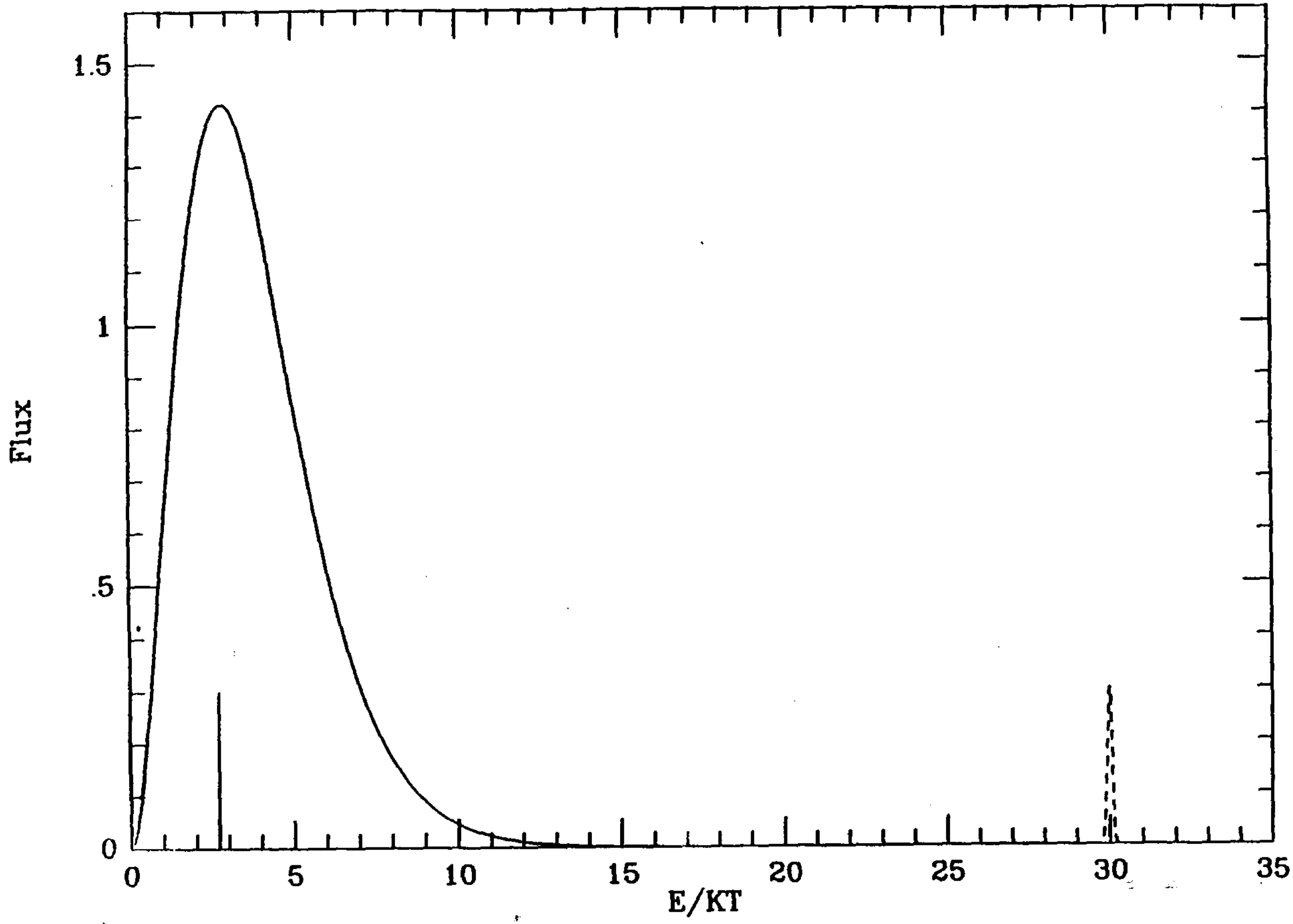
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