Formation of Galaxies in a Gravitino-Dominated Universe

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If gravitinos of mass 1 keV (or similar particles) dominate the mass of the universe, a critical scale of galactic size arises due to their collisionless phase mixing. It is shown that density perturbation spectrum of gravitinos is relatively flat between galactic and cluster scales, unlike the massive neutrino case. If gravitinos form the dark matter, initial adiabatic fluctuations lead to a hierarchical picture of clustering. Galaxies form first, but dissipation is necessary for their survival.

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One of the crucial questions facing astrophysics is the nature of the dark matter which dominates the mass density of the universe. No theory of the formation of large scale structure can be complete without understanding its role in clustering. From dynamical studies of galaxy clusters,[1] Ω, the ratio of the average to critical mass density of the universe, is estimated to be of order unity, yet the luminous matter in galaxies accounts for only about one percent of that amount. It seems unlikely that the dark matter could have been baryonic at the time of big-bang nucleosynthesis, since the observed He and D abundances require[2] $\Omega_B \leq 0.1$, making pregalactic black holes, low-mass stars, and hot intergalactic gas unattractive as explanations of all the dark matter. The ideal candidate is nonbaryonic, and is probably nondissipative since dark matter appears to be less condensed than the luminous component.[3] Massive, weakly interacting, stable neutral particles satisfy these requirements. One attractive possibility is a neutrino with mass 10–100 eV.[4] Such a neutrino would dominate the mass density of the universe ($\Omega_\nu \geq \sum m_{\nu i}/100$ eV), and modify the usual theories of galaxy formation. After neutrino decoupling at $T \approx 1$ MeV, they are collisionless, and hence are subject to phase mixing. As a consequence, primordial adiabatic perturbations in the neutrino sea on mass scales less than a critical value are strongly damped, resulting in a fluctuation spectrum which is sharply cutoff below the mass scale[5,6] $M_{\nu m} \cong 1.8 m_P{}^3 m_\nu{}^{-2}$ where m_P is the Planck mass. Since $\Omega_{tot} \leq 2$, $m_\nu \leq 200$ eV, and $M_{\nu m} \geq 10^{14} M_\odot \gg M_{gal} \approx 10^{12} M_\odot$ superclusters of mass $(10^{15}-10^{16}) M_\odot$ form first.

Here we assume the universe to be dominated by a stable, more weakly interacting neutral particle (hereafter X), e.g., a gravitino,[7] photino,[8] or a right-handed neutrino.[9] Neutrinos decouple just before $e^\pm$ pairs annihilate; thus today the temperature of the ν momentum distribution is less than that of the photons, $T_\nu = (4/11)^{1/3} T_\gamma$. Because the X interacts more weakly, it decouples much earlier and does not share in the heating caused by the annihilation of the various species that occurs after X decoupling (e.g., $\mu^\pm$, $\pi^\pm\pi^0, e^\pm$). Today X's are cooler and less abundant than ν's. If X's decouple when they are relativistic, their abundance relative to photons is $n_X/n_\gamma = \frac{1}{2} g_X (T_X/T_\gamma)^3 = 1.9 g_X/g_*$, where $T_X = (3.9/g_*)^{1/3} T_\gamma$ is the temperature of the X momentum distribution, g_X is N_X (the number of degrees of freedom) for a boson, $3N_X/4$ for a fermion, and $g_* = \sum N_B + \frac{7}{8}\sum N_F$ is the total effective number of degrees of freedom of all relativistic species in thermal equilibrium at X decoupling. If $m_X \geq (30\ \text{eV})(g_*/100) \times (g_X/1.5)^{-1}$, the X dominates the mass density of the universe ($\Omega_X \geq \Omega_B$) and since $\Omega_{tot} \leq 2$, it must be lighter than[9] $m_X \leq (2\ \text{keV})(g_*/100)(g_X/1.5)^{-1}$. Pagels and Primack[7] have suggested that the $\pm\frac{1}{2}$ helicity states of the gravitino, the supersymmetric partner of the graviton, has a mass ≈ 1 keV and decouples at $T = 100$ GeV when $g_* = 100$, making it a candidate for the dark matter. We shall use a particle with these parameters as our generic example; however, we shall exhibit how our results scale with m_X, g_*, and g_X. In a ν-dominated universe $n_\nu \approx n_\gamma$, neutrinos dominate the energy density shortly after they become nonrelativistic (NR). In an X-dominated universe $n_X/n_\gamma \ll 1$, so the X does not dominate the energy density until well after it becomes NR. This affects the evolution of density perturbations, as we now show.

Various mass scales can be defined which characterize the perturbation spectrum (Fig. 1). These are all of the same order of magnitude for ν's and lead to a steep cutoff: the characteristic scale is fairly independent of the initial spectrum. For the gravitinos these scales are quite different, and result in a spectral peak which cannot simply be scaled to the value $M_{Xm} = 1.8 m_P^3 m_X^{-2}$; and as we will show, the spectral shape is very sensitive to the initial spectrum.

Primordial linear density perturbations in the X's have a density contrast $\delta = \delta\rho/\rho$ which depends upon the mass scale of the perturbation M, a measure of its (comoving) size.[10] The gravitational attraction of an overdense region in the X's balances the dispersive effects associated with thermal motion at the Jeans mass given by $M_J = 42 m_P^3 m_X^{-2} g_*^{-2} g_X (a/a_{eq})^{-3/2}$ in the NR regime. Here a is the scale factor of the universe, normalized to unity at $T_X = m_X/3.15$; $a_{eq} = 146(g_*/100)^{4/3}(g_X/1.5)^{-1}$ is its value when the energy density in X's equals that of the photons and neutrinos. For the massive ν's, damping

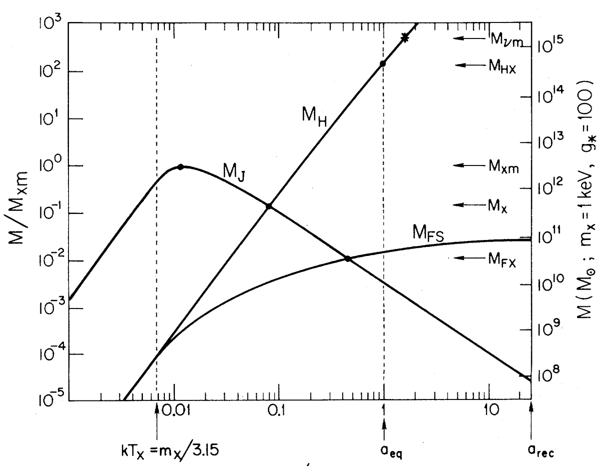

FIG. 1. The Jeans mass (M_J), horizon mass (M_H), and free-streaming mass (M_{FS}) for a gravitino-dominated universe are plotted against scale factor. The mass scales characterizing the perturbation spectrum are also shown.

occurs for scales below the peak in their Jeans mass $M_{\nu m}$. This is not the case for X's: it is necessary for them to have propagated over the scale of the perturbation before significant damping can occur. If they move at the speed of light, they can only have covered the scale of the horizon, $M_H = 0.9 m_P^3 m_X^{-2} g_*^2 g_X^{-2}[(1 + a/a_{eq})^{1/2} - 1]^3$. Certainly damping cannot occur until $M_H > M_J$; thus a characteristic scale is set when $M_J = M_H$:

$$M_X = 11.6 m_P^3 m_X^{-2} g_*^{-2/3}. \quad (1)$$

In fact, this is an overestimate of the damping mass, since the X's are NR, and slow down as the universe expands. If the X's freely stream, they cover a region defined by the mass scale $M_{FS} = 8.6 m_P^3 m_X^{-2} g_*^{-2} g_X (1 + \ln a)^3$, valid for $a \leq a_{eq}$. Perturbations on scales smaller than both M_J and M_{FS} are expected to be erased, while those in excess of either M_J or M_{FS} cannot have damped very much. This defines a scale M_{FX} found by equating M_J and M_{FS}:

$$M_{FX} = 0.44 m_P^3 m_X^{-2} g_*^{-4/3} g_x. \quad (2)$$

The final spectral shape is determined not only by the damping but also by the details of fluctuation growth. If the perturbations are adiabatic, which seems to be the most likely initial condition, then the fluctuations will be of equal magnitude in the X's, photons, and ν's, and they grow together until the horizon has expanded to encompass them. Upon entering the horizon, perturbations in the massless ν's decay away and those in the photons oscillate at constant amplitude. The photons, tightly coupled to the baryons, behave as an ideal fluid prior to the recombination of hydrogen when $T_\gamma = 3000$ K. The oscillation period, proportional to the light crossing time across the perturbation, is generally much shorter than the response time for NR X's. Therefore both the radiation and the ν fluids appear unperturbed as far as the X's are concerned and cannot drive fluctuation growth in the X gas for $M < M_H$. The behavior of the perturbation after it enters the horizon is described by[11] $\delta_X \approx 1 + \frac{3}{2} a/a_{eq}$, implying little growth for $a \leq a_{eq}$. This retardation does not occur once X's dominate the energy density; perturbations which come within the horizon after a_{eq} have uninterrupted growth, so the horizon mass at a_{eq} sets another scale:

$$M_{HX} = 0.9 m_P^3 m_X^{-2} g_*^2 g_X^{-2}. \quad (3)$$

The horizon mass, Jeans mass, and free-streaming mass are plotted in terms of the scale factor in Fig. 1. These are the characteristic scales for the X perturbations. Silk damping[12] due to the viscous coupling of baryons to γ's near recombination erases baryon perturbations on scales below[6] $M_s \sim 10^{16} M_\odot$ for $\Omega_B \sim 0.01$. However, baryons fall into the gravitino-generated potential wells, and by a red shift, $z \sim 100$, the baryon perturbation spectrum will match that of gravitinos, and the two will evolve toward nonlinearity together. The perturbations grow by a factor $a_{\rm now}/a_{\rm eq} \sim 28\,000$ by the present epoch at $a_{\rm now} = 4\times10^6(g_*/100)^{1/3}(m_X/1 \text{ keV})$ for $M < M_{HX}$; this is similar to the growth factor in a ν-dominated universe. To recapitulate, we expect damping below M_{FX} and retarded growth below M_{HX}. If initially the spectrum is a power law, $\delta\rho/\rho \sim M^{-n/6-1/2}$, then we expect the final spectrum to maintain its initial shape for $M > M_{HX}$, and to be flattened by a factor of $M^{2/3}$ for $M_{HX} > M > M_X$, reflecting the values of $\delta\rho/\rho$ on the horizon in this range. Some effects of damping will be felt between M_X and M_{FX}, and there will be a rapid falloff below M_{FX}.

To compute the final X perturbation spectrum, we solved the linearized Einstein equations describing a universe perturbed about an expanding Friedman-Robertson-Walker background, coupled to the linearized collisionless Boltzmann equation for the X's and ν's and to hydrodynamic equations for the γ's. This system of coupled Einstein-Boltzmann equations was solved numerically in the same manner as for a ν-dominated universe,[13] with only slight modifications needed to treat the X-dominated case. Peebles[14] has done a similar calculation recently. The numerical integration for the evolution of each spatial Fourier component of the metric and density coefficients was started when the horizon was much smaller than the wavelength and was terminated only when the asymptotic growing mode in the NR regime had been reached. The solutions, though calculated from a flat "white-noise" spectrum ($n=0$), can be scaled to an arbitrary initial spectrum by multiplying with the appropriate powers of k ($|\delta_k|^2{}_i \sim k^n$). We plot the resulting $\delta\rho/\rho(M) \sim k^{3/2}|\delta_k|$ in Fig. 2 for the following initial spectra: white noise ($n=0$), constant curvature ($n=1$), causal limit ($n=4$). Here $M\sim k^{-3}$ relates the comoving wave vector, k, to the mass scale,[10] and $\delta\rho/\rho(M)$ measures the average relative contribution of the scale M to local density enhancements. The spectra for a ν-dominated universe are also plotted for comparison. These show a

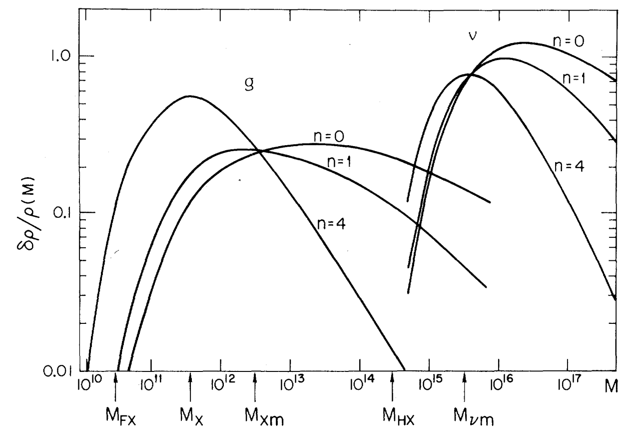

FIG. 2. The final density perturbation spectrum of gravitinos (g) for different initial power law spectra characterized by the index n is plotted against mass scale. For comparison the spectra arising in a neutrino-dominated universe (ν) are also shown. The normalizations of all curves are arbitrary.

much steeper falloff from the peak not only to the small-mass end due to the strong damping, but also to the large-mass end due to the change in the expansion rate after $a = a_{\rm eq}$. Still, even for X perturbations, there is not much structure below M_{FX}, and damping is already effective in slowing growth below M_X.

The interpretation of the $n=4$ gravitino spectrum is straightforward. The first scale to reach nonlinearity $\delta\rho/\rho \sim 1$, and begin to collapse, is $M_X \approx 3\times10^{11} M_\odot$, which corresponds to galactic scales. At later times larger mass scales reach nonlinearity and collapse, resulting in hierarchical clustering of galaxies. The $n=1$ spectrum will lead to similar behavior, though collapse on larger scales will occur sooner than in the $n=4$ case.[15] The $n=0$ spectrum is quite flat over the range $(10^{12}$ to $10^{15})M_\odot$, and perturbations on these scales go nonlinear at approximately the same time. If $n \leq -1$, the peak is at $\sim 5\times10^{14} M_\odot$: clusters collapse first; substantial substructure may form due to fragmentation.

In order for galactic-scale structures to survive during the collapse of larger systems, it is necessary to invoke dissipation on time scales less than freefall, since White and Rees[16] have shown that substructures formed by collisionless processes would be destroyed. Numerical N-body simulations[17] support this view. In any case, dissipation is required in galaxies to account for the separation of the luminous from the dark matter. Galaxies in groups and clusters will then have their collisionless halos partially or totally tidally stripped, and yet their luminous cores will survive. The mass of baryons in a galaxy of total mass M_X is only $(\Omega_B/\Omega_X)M_X \sim 10^{10} M_\odot$, and not all of this may be able to cool sufficiently to avoid being stripped along with the loosely bound gravitinos. Low-velocity gravitinos should be left behind. Since it is more difficult to eject gravitinos from deeper potential wells, one expects larger mass objects to have kept more dark matter than smaller ones, resulting in a mass-to-light ratio increasing with scale, a result which is inferred from recent red-shift surveys.[3] Though M_X defines a galactic scale, the ability of the central baryons to cool in a Hubble time is necessary for the formation of galaxies to occur, and indeed sets a mass scale as well.

In the X-dominated universe, superclusters would be identified with the last scale to have collapsed, perhaps $M_{HX} \sim 10^{15} M_\odot$. Holes in the galaxy distribution would also arise on these scales, evacuated by infall into superclusters. The infalling objects (galaxies and clusters) would presumably be nondissipative by this stage; this could lead to high velocities. Superclusters are observed to be flat,[18] with low peculiar velocities, which suggests dissipation upon formation. These observations favor a ν-dominated universe, in which superclusters and large holes in the galaxy distribution naturally appear on scales $M_{\nu m}$.

We have presented computations of the density fluctuation spectrum in a gravitino-dominated universe. Galaxy scales are naturally expected to arise. Even though the theory is an adiabatic one, clustering is generated in a hierarchical manner with galaxies forming clusters, then superclusters as in the White and Rees scenario. This contrasts with the neutrino-dominated adiabatic picture in which flat superclusters form first, and then galaxies arise from pancake fragmentation, as Zeldovich[19] originally suggested. The identification of the dark matter with massive neutrinos or gravitinos therefore leads to theories of galaxy formation which are similar to the original adiabatic and isothermal theories, respectively. Neither can be ruled out at the present time, but the observational data on large scale clustering have undergone dramatic improvements over the last years and will ultimately decide which theory is preferred, and thus shed light on the nature of the dark matter. Particle physics theories which predict stable, weakly interacting, neutral particles must then be confronted with what is known about galaxy formation.

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