SATELLITE SURVIVAL IN COLD DARK MATTER COSMOLOGY

C. M. Boily
Observatoire astronomique, 11 rue de l’Université, F-67000 Strasbourg, France; cmb@astro.u-strasbg.fr

N. Nakasato
Department of Astronomy, University of Tokyo, Bunkyo-ku, Tokyo 113-0033, Japan

AND

R. Spurzem and T. Tsuchiya
Astronomisches Rechen-Institut, Mönchhofstrasse 12-14, D-69120 Heidelberg, Germany

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ABSTRACT

We study the survival of substructures (clumps) within larger self-gravitating dark matter halos. Building on scaling relations obtained from N-body calculations of violent relaxation, we argue that the tidal field of galaxies and halos can only destroy substructures if spherical symmetry is imposed at formation. We explore other mechanisms that mayetail the number of halo substructures during the course of virialization. Unless the larger halo is built up from a few large clumps, we find that clump-clump encounters are unlikely to homogenize the halo on a dynamical timescale. Phase mixing would proceed faster in the inner parts and allow for the secular evolution of a stellar disk.

Subject headings: cosmology: theory — dark matter — galaxies: dwarf — methods: N-body simulations

1. INTRODUCTION

High-resolution simulations of structure formation in a ΛCDM cosmogony have revealed a large number of substructures within dark matter halos (e.g., Ghigna et al. 2000; Moore 2001). The mass function of these dark clumps is (Moore et al. 1999b; De Lucia et al. 2004)

\[ n(m) \propto m^{-2} \]

(1)

in the mass range \(10^8 - 10^{11} M_\odot\). These orbiting self-bound dark matter satellites would drag along baryonic matter to a nonnegligible fraction of their mass. Consequently galaxies should harbor a large number of dwarf galaxies, while only a handful are found (Kauffmann et al. 1993; Moore et al. 1999b; see Binney & Merrifield 1998).

The situation is made worse from a dynamical standpoint. Massive dark clumps would perturb the vertical structure of galactic disks through tidal heating, while their thin structure suggests that the immediate neighborhood of disks is devoid of such perturbations (Toth & Ostriker 1992; Moore et al. 1999b). The large-mass end of the clump mass distribution function is robust in view of the fact that simulation checks find it to be robust to numerical resolution issues (Moore et al. 1999a; Gao et al. 2004; Power et al. 2003). Therefore, the orbital distribution of clumps in phase space must allow for long periods of unperturbed evolution by the disk, as emphasized by Navarro (2002) and Font et al. (2003). A more severe problem is the narrowness of streams of stars associated with dwarf galaxies’ tidal debris (Ibata et al. 2001; Johnston et al. 2001), which suggests a smooth background halo potential in order to preserve the cohesion of the stream, in direct conflict with computer simulations of structure formation. Observational evidence drawn from solar neighborhood kinematics points to a fine meshing of low-mass streams to account for coherent motion into and out of the galactic plane (Helmi et al. 2002; Gould 2003). Settling this issue requires both a precise map of the morphology of halos and convergence in the mass distribution function of clumps with a resolution down to the subdwarf mass range, still a challenge to present-day computer models (see, e.g., Power et al. 2003). The morphology of halos has been discussed at length by several authors (e.g., Moore 2001; Ghigna et al. 2000; Fukushige & Makino 2003; Power et al. 2003; Navarro et al. 2004). Several contributions in Natarajan (2002) give a broad overview of the formation processes and equilibrium properties of halos derived from numerical simulations and possible observational tests for their detection.

Particle-based calculations of galaxy formation proceed with a number routinely approaching a few \(10^7\) particles for a whole simulation. The mass of individual particles that make up dark halos in these simulations, \(~10^6 M_\odot\), is still large compared with a mass spectrum of dark matter that may yet include a population of brown dwarfs (i.e., stellar masses). Furthermore, on the order of \(~10^3\) mass elements must participate in the formation of a self-gravitating body to resolve the growth of potential energy adequately (Boily et al. 2002; Roy & Perez 2004), or a mass resolution of \(10^8/10^5 \approx 10^3 M_\odot\) is required to account for dwarf-size structures. This leaves a gap in mass resolution of some 3 orders of magnitude from that achieved by present-day simulations. In this context, we need to establish when the statistics of dark matter clumps derived from N-body computations can be scaled up to actual galactic systems. A crude picture of halo formation divides the process into two stages: one of rapid collapse on a free-fall timescale, followed by a second, longer period of sporadic accretion (Bower 1991; Lacey & Cole 1994; Zhao et al. 2003a, 2003b). The first stage involves structures spread over a narrow range of mass, while the second stage sees low-rate accretion by a

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1 Corresponding author.

2 Current address: Institute of Advanced Physical & Chemical Research (RIKEN), 2-1 Hirosawa, Wako-shi, Saitama 351-019, Japan.

3 Now at SGI Japan, Ltd., Ebisu Garden Place Tower 31F, 4-20-3 Ebisu, Shibuya-ku, Tokyo 150-6031, Japan.
dominating, central body. The growth of potential energy during the free-fall stage of formation sets the maximum phase-space density of the distribution function by efficiently, if incompletely, redistributing binding energy between particles (Lynden-Bell 1967; van Albada 1982).

The growth of potential energy during collapse of self-gravitating systems is a function of the initial morphology and the number of particles (or equivalently, mass resolution) used in the computation (e.g., Boily et al. 2002 and references therein). A direct consequence of this is the maximum tidal field, $\propto \nabla^2 \phi(r)$, experienced by galactic satellites (dark or baryonic) as they fly through the system varies with the morphology of the underlying halo, i.e., its formation history. For instance, at fixed resolution, the strength of tidal fields reaches a higher maximum for an initially spherical distribution than for an axisymmetric or triaxial one. The simplification of symmetric distributions allows the scaling up of the results with particle number exactly (see §2 below). Would the tidal heating experienced by dark clumps during infall, scaled up to actual galactic halo particle numbers, be sufficient to unbind them? In this short contribution, we apply the results of violent relaxation studies to the tidal heating of galactic substructures in a hierarchical Einstein–de Sitter universe. We show that the tidal field developing during infall may yet be sufficient to erode substructures smaller than a critical linear size $l_c$ if the particle number $N$ is sufficiently large and the galactic halo shows axial or spherical symmetry during infall.

2. SCALING OF TIDAL FIELDS

Aarseth et al. (1988) have shown that the growth rate of global modes of fragmentation during violent relaxation is such that the minimal radius achieved by a spherical distribution scales with the simulation particle number $N$ as

$$ C \equiv \frac{R(0)}{R(t_f)} \propto N^{1/3}, \quad (2) $$

where $R(0)$ is the initial system size enclosing $N$ identical mass elements, and $R(t_f)$ is the radial size at the free-fall time $t_f$ defined by

$$ t_f = \sqrt{\frac{3\pi}{32G\langle\rho(a,0)\rangle}}, \quad (3) $$

where $\langle\rho(a,0)\rangle = 4\pi M/3a^3$ is the mean density inside the particle’s initial radius $a$. A recent study extended result (2) empirically with $N$-body calculations to axisymmetric (cylindrical) and triaxial initial configurations (Boily et al. 2002). It was found in these cases that $C$ scales with particle number as

$$ C \propto \begin{cases} N^{1/6}, & \text{axisymmetric;} \\ \text{constant}, & \text{triaxial.} \end{cases} \quad (4) $$

The factor $C$ achieved by systems starting from triaxial distributions remains $\lesssim 40$ for $\sim 86\%$ of the parameter space of axial ratios $a : b : c$; however, the maximum achieved by individual realizations is highly sensitive to the initial axes’ ratios and may yet diverge. $N$-body simulations using direct-integration algorithms (Aarseth et al. 1988; Theis & Spurzem 1999; Boily et al. 1999) or fast Fourier transform integrators (e.g., Boily et al. 2002) reproduce these relations over several decades in particle number, giving confidence that the results are well recovered independently of the numerical method used. A consequence of equations (2) and (4) is that a spherical or axisymmetric distribution will contract ever more as the number of particles $N$ is increased. The morphological evolution of halos and their substructures will not be well resolved if at formation the particle number that takes part in the relaxation phase is too low. Zhao et al. (2003a, 2003b) show that the formation history runs through a rapid accretion phase followed by slow growth of the outer halo. The scaling laws (2) and (4) obtained for violent relaxation would therefore apply best to the first, early phase of formation, but not the late one.

3. COSMOLOGICAL APPLICATION

The mean tidal field of a collapsing halo is evaluated from the time-dependent second derivative of the gravitational potential by

$$ F_t = \nabla^2 \Phi|_R = -\frac{GM}{R^3} \quad (5) $$

with $R$, and double-differencing at fixed time yields a measure of the tidal field at $R$ acting on a clump of size $l_c$. It is clear that the tidal force is unbound if during collapse $R(t_f) \rightarrow 0$, as would occur, e.g., for a large-$N$ spherical distribution (cf. eq. [2]). Bound clumps would survive violent relaxation if their binding energy were higher than the tidal heating that they suffer during that phase (van Albada 1982; Tsuchiya 1998). As shown from equation (5), remnant structures will be severely disrupted if the maximal potential depth achieved is large. To progress further we need to invoke a result for structure formation in an expanding universe. Present-day data support an asymptotically flat metric for the universe, and hence the Einstein–de Sitter cosmogony remains attractive.

3.1. Background Tidal Heating

In an Einstein–de Sitter universe the relation between a bound structure’s mass and virial radius is (Kaiser 1986; Padmanabhan 1993; Somerville & Primack 1999)

$$ R \propto M^\gamma, \quad (6) $$

where $\gamma$ is known in terms of the power spectrum of density fluctuations $P(k)$ at wavenumber $k$. The classic CDM power spectrum at the time $\tau$ of galaxy formation accounts for the time evolution of structure from a bottom-up point of view. The relation between $P(k; \tau)$ and the Zel’dovich spectrum arising from postinflation decoupling $P(k; t_i)$ is (see Bardeen et al. 1986, Appendix G)

$$ P(k; \tau) = \left[\frac{a(t_i)}{a(\tau)}\right]^2 T^2(k)P(k; t_i), \quad (7) $$

where $a$ is the cosmic expansion factor and $T$ is the transfer function, which is well approximated analytically by

$$ T(q) = \frac{\ln (1 + 2.34q)}{2.34q} \times (1 + 3.9q + 259q^2 + 163q^3 + 2027q^4)^{-1/4}, \quad (8) $$
where \( q \equiv |k[\theta^{1/2}/(\Omega h^2 \text{Mpc}^{-1})] \), and \( \theta = \rho_{\text{rel}}/1.68\rho_c \), is the ratio of relativistic particles to photon energy densities. On the largest scales the Zel’dovich spectrum \( P(k; \tau) \propto P(k; t_i) \propto k \) is recovered from equation (8), while on small scales \( P(k; \tau) \propto k^{-(n_k+3)} \). If we fit the power spectrum locally to a power law of index \( n_k \), \( P(k) \propto k^{n_k} \), the index shifts progressively from \( n_k = +1 \) to \( n_k = -3 \) as we explore smaller scales. The indices \( n_k \) and \( \gamma \) are linked through mass-radius relation (6) by

\[
\gamma = (n_k + 5)/6.
\]

Note that \( \gamma \) is related to the power index \( \nu \) relating mean density and mass, \( \rho \propto M^\nu \), by \( \nu = \gamma - 3 = (n_k - 13)/6 \); the free-fall time \( t_{\text{ff}} \propto 1/\rho^{1/2} \) therefore scales with mass as

\[
t_{\text{ff}} \propto M^{(13-n_k)/12}.
\]

On the smallest scales \( n_k \to -3 \) and \( t_{\text{ff}} \propto M^{4/3} \), a steeper relation than on the largest scales when \( n_k \to +1 \) and \( t_{\text{ff}} \propto M \). Therefore small clumps have fully virialized when the violent relaxation phase of the larger halo begins.

Since there are no fixed scales of mass or radius in gravitational dynamics, all virialized structures obey the same relation for equation (6). If we lump together all those of virial radius less than \( r \), of mass \( m \) chosen such that \( M/m = N \gg 1 \), we have

\[
\frac{R}{r} \propto \left( \frac{M}{m} \right)_{\gamma} = N^{(n_k+5)/6}.
\]

In a hierarchical universe, small structures form first and hence the \( N \) small clumps have virialized well before the large underlying halo of total mass \( M \). At constant mass the virial theorem provides a relation between equilibrium radius and initial size \( R_i \) for a self-gravitating system,

\[
R_i = 2R,
\]

which applies equally to all structures. We need to relate the tidal field (5) with equations (2) and (4) to determine whether a structure of size \( r \) survives the formation of the larger halo of radius \( R_o \). The energy transferred to a small clump by tidal forces during infall is adequately quantified by Spitzer’s (1958) impulse approximation formula even for relatively slow encounters (Aguilar & White 1985). This gives confidence that it will hold in the present context, where velocities are high. Tsuchiya (1998) finds this to be correct in his study of relaxing Plummer distributions. The tidal energy \( \Delta E \) gained by a substructure of mass \( m \), radius \( r \), and internal binding energy \( E \) can be evaluated for a single passage at velocity \( V \) across the background halo (or galaxy) potential at the time of collapse, when the tide is maximum. In the impulse approximation, this is given by

\[
\Delta E = \frac{4G^2M^2m}{5R(t_f)^3V^2} \approx \left( \frac{\bar{r}}{R_i} \right)^3 C_3 \frac{M}{m} E.
\]

Note that equation (11) applies to the tidal field of the background potential, and does not account for individual clump encounters. These are discussed later. Substituting equation (9) in equation (11) we get

\[
\Delta E = N^{-(n_k+3)/2} \left( C_3 \right)^{3/2} E, \quad (12)
\]

whence we deduce that \( \Delta E \ll E \) if the coefficient \( N^{-(n_k+3)/2}E \) remains small. For spherical systems the collapse factor \( C \) obeys equation (2) and therefore

\[
\frac{\Delta E}{E} \propto N^{-(n_k+1)/2}, \quad \text{spherical},
\]

while for axisymmetric or triaxial distributions we find from equation (4)

\[
\frac{\Delta E}{E} \propto \begin{cases} 
N^{-(n_k+2)/2}, & \text{axisymmetric;} \\
N^{-(n_k+3)/2}, & \text{triaxial.}
\end{cases} \quad (13)
\]

The consequences of these results in relation to the power spectrum of density fluctuations are clear: for clumps orbiting in a collapsing spherical halo or galaxy, tidal heating will be ineffective provided \( n_k > -1 \). If the underlying distribution is axisymmetric, tidal heating will be ineffective when \( n_k > -2 \). However, for triaxial initial conditions this will hold true if \( n_k > -3 \). Since the power spectrum of the observed matter distribution is never steeper than \( n_k = -3 \), we deduce that the bulk of substructures or clumps evolving in larger structures, such as dark matter halos, will survive the formation of triaxial distributions, if gravity alone fixes their binding energy.

### 3.2. Tidal Heating Due to Other Substructures

The above conclusion only concerns the response of clumps to the background tidal field. We can consider the interaction between clumps themselves as they cross the dense system. To this end we consider the tidal heating by two equal-mass substructures during an encounter. Substituting \( M \to m \) in equation (11) and the radius \( R(t_f) \) by the mean distance between clumps \( x \approx R(t_f)/N^{1/3} \), an effective impact parameter, we find

\[
\Delta E = \frac{4G^2m^3}{3(R/N^{1/3})^3V^2} \bar{r}^2 \approx 2m \frac{2}{3} M N^{4/3} \left( \frac{\bar{r}}{R} \right)^3 E.
\]

In the above we used \( V^2 \approx 2GM/R \), with \( M = Nm \) the total system mass as before. The expression reduces to

\[
\Delta E \approx \frac{2}{3} N^{1/3} \left( \frac{\bar{r}}{R} \right)^3 E. \quad (14)
\]

Clearly, to achieve \( \Delta E \approx E \) requires \( r \sim R \) or \( N \gg 1 \). We can simplify equation (14) by substituting for \( \bar{r} \) using equation (9). We then find

\[
\Delta E \ll \frac{2}{3} N^{-(13+3N)/6} E. \quad (15)
\]

Thus for any appreciable number \( N \) the tidal heating due to encounters between clumps will be significant if \( n_k < -13/3 \approx -4.3 \). No regime of the power spectrum covers that range and hence encounters between clumps never produce significant tidal heating. A direct consequence of this is that while halos form at different redshifts and sample different regimes of
the structure power spectrum, the mass distribution function of substructures should be robust against cutoffs or significant changes to its shape. It is not clear yet whether the scale-free nature of the clump mass function (eq. [1]) measured in N-body simulations (e.g., Ghigna et al 2000) can be extended to very small masses (see Gao et al. 2004).

There is, of course, one situation in which substructures can heat one another through tidal forces, which is when \( r \sim R(0) \) or \( R(t_{b}) \) and \( N \sim 1 \). Indeed, when \( r \) matches the mean interclump distance \( \sim R(t_{b})/N^{1/3} \), we compute \( \Delta E/E \sim 1 \), always. However, this situation is more appropriate to galactic mergers than the formation of a halo through accretion of several subunits, as would occur in any bottom-up calculation of galaxy formation.

The survival of substructures is in part due to the large relative velocity \( V \) established under the mutual potential of all clumps. Survival during infall is no guarantee that the substructures would remain once the halo has achieved virial equilibrium. Estimating more precisely the net rate of heating on one clump due to the background tidal field as it crosses the system is made difficult because of the large changes in potential taking place during violent relaxation. This is best done with numerical N-body calculations tailored for this problem. The very high resolution simulations performed by, e.g., Ghigna et al. (2000) demonstrate the likely survival of most substructures in and around dark matter halos postvirialization, in support of the basic argument outlined here.

4. DISCUSSION AND CONCLUSION

Galactic satellites survive the formation phase of triaxial halos and galaxies and will be destroyed on long timescales as they orbit the host galaxy. This result is obtained both from large N-body simulations (see also, e.g., Bullock et al. 2000; Ghigna et al. 2000) and from semianalytical arguments (Moore et al. 1996; Taffoni et al. 2003) as well as the fluid calculation presented here, and is therefore robust. The initial conditions and subsequent evolution of N-body computer simulations still plague their interpretation and application to observed galaxies in terms of simple estimates of satellite disruption times. For instance, Font et al. (2003) and Ardi et al. (2003) have questioned the rate of disk heating by infalling dark satellites. These authors find that thin disks may yet remain stable despite a high count of bound dark matter clumps, provided the clumps do not follow near-radial orbits. The problem of disk heating would seemingly not occur if the inner region of the halo were isotropic in phase space. We already noted that the destruction of dark satellites would be more effective in the deep potential of the inner halo. Recently, Gao et al. (2004) have rederived statistics of halo substructures in computer simulations and found them to be less concentrated than the host halo. This and the coherence of cold-stream satellite debris could be interpreted in the light of the present analysis as pointing to a near-spherical (and hence destructive), early phase of halo formation. Several studies have argued for a more spherical morphology in the inner region of halos (e.g., Blumenthal et al. 1986; Dubinsky 1994). The cooling of baryons at the heart of DM halos would provide a mechanism for this by locking the inner halo morphology to a rounder shape than that obtained from strictly gravitational evolution (Dubinsky 1994; Frenk et al. 1996). The inner morphology of galactic halos would not automatically be spherical if baryons have had time to cool and form disks before the halo assembly is complete; when that is the case, the halo’s inner morphology in equilibrium is even more sensitive to the formation history and depends, for instance, on the orientation of the disks as they merge (see Kazantzidis et al. 2004).

Possible mechanisms that may disrupt dark clumps on the dwarf galaxy mass scale and below include supernova blowouts through gas irradiation and expulsion (Efstathiou 1992; Somerville 2002; Gnedin & Zhao 2002). The net effect of gas loss on gravitationally bound structures is unlikely to be significant if the gas mass fraction is small (Hills 1980; Boily & Kroupa 2003). Gnedin & Zhao (2002) have argued that the peaked density profiles obtained from CDM numerical calculations would resist rapid removal of the gas under any realistic circumstances. Thus, unless the dwarf-size clumps contain a very large fraction of baryonic matter, they will survive any degree of gas heating.

Côté et al. (2002) presented a Monte Carlo simulation of chemical enrichment of galaxies through gas evaporation from clumps of dark matter initially seeded with baryons (uniform \( M/L \) ratio). They find that the mass function of seeded clumps required to match their sample of galaxies (in terms of chemical gradients and observable dwarf galaxy and star cluster populations) is similar to the mass function (1) obtained from large-N cosmological calculations. This would suggest that the high-mass end of the clump mass function survives the formation of the host dark halo to produce the observed population of dwarfs. It does not, however, suppress the number of dark clumps that may still be orbiting the halo. Another route to solving the overabundance of dwarf galaxies is preventing (bright) baryons from forming stars. This can be achieved either through background UV radiation (Efstathiou 1992; Somerville 2002) or by preventing a Toomre instability from developing fully (Verde et al. 2002), effectively shutting off the formation of stars in the first place. Gas-rich dwarfs would undergo substantial morphological evolution through ram stripping from the intergalactic medium; their long-term fate (destruction or survival) must account for such evolution, since it will change the dwarf binding energy through dissipation (Mayer et al. 2001).

Dynamical friction can in principle provide an alternative solution if the clumps spiral in rapidly and lose mass owing to tidal heating (Syer & White 1998; Tormen et al. 1998). Computer simulations and analysis of decaying satellites show that a heavy satellite loses up to 90% of its mass in a few orbital periods (e.g., Klessen & Kroupa 1998; Peñarrubia et al. 2002; Taffoni et al. 2003; see also Hashimoto et al. 2003). Van Kampen (2002) has argued that the effect of dynamical friction may yet be underestimated in computer simulations because of limited mass resolution. Dynamical friction, in conjunction with tidal forces, will cause the disruption of a satellite after a period of time (e.g., Ibata et al. 1994; Klessen & Kroupa 1998; Bullock et al. 2001). Bullock et al. (2001) argue that the halo stellar population of the Milky Way may be accounted for if a sufficient number of satellite dwarf galaxies had already accreted their mass at the time the galactic halo formed and were then stripped of their less-bound stars by galactic tides. Clearly the links between halo morphology, halo substructure statistics, and stellar populations offer more avenues for future work.

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