Interaction of clumpy dark matter with interstellar medium in astrophysical systems.

A. N. Baushev
Bogoliubov Laboratory of Theoretical Physics, Joint Institute for Nuclear Research, 141980 Dubna, Moscow Region, Russia
DESY, 15738 Zeuthen, Germany
Institut für Physik und Astronomie, Universität Potsdam, 14476 Potsdam-Golm, Germany

ABSTRACT
Contemporary cosmological conceptions suggest that the dark matter in haloes of galaxies and galaxy clusters has most likely a clumpy structure. If a stream of gas penetrates through it, a small-scale gravitational field created by the clumps disturbs the flow resulting in momentum exchange between the stream and the dark matter. In this article, we perform an analysis of this effect, based on the hierarchial halo model of the dark matter structure and Navarro-Frenk-White density profiles. We consider the clumps of various masses, from the smallest up to the highest ones $M \geq 10^9 M_\odot$. It has been found that in any event the effect grows with the mass of the clump: not only the drag force $\vec F$ acting on the clump, but also its acceleration $\vec w = \vec F / M$ increases.

We discuss various astrophysical systems. The mechanism proved to be ineffective in the case of galaxy or galaxy cluster collisions. On the other hand, it played an important role during the process of galaxy formation. As a result, the dark matter should have formed a more compact, oblate, and faster rotating substructure in the halo of our Galaxy. We have shown that this thick disk should be more clumpy than the halo. This fact is very important for the indirect detection experiments since it is the clumps that give the main contribution to the annihilation signal. Our calculations show that the mechanism of momentum exchange between the dark and baryon matter is ineffective on the outskirts of the galactic halo. It means that the clumps from there were not transported to the thick disk, and this regions should be more clumpy than the halo on the average.

Key words: cosmology: dark matter, cosmology: theory, elementary particles.

1 INTRODUCTION

In modern representation about 80% of the matter forming structures in the Universe is dark matter (DM). In particular, it is the dark matter that makes the main contribution to the contents of galactic haloes. The nature of the dark matter is presently unknown. It is widely believed that it consists of some weakly or extremely weakly interacting particles generated in the early Universe. Modern cosmological observations (Spergel et al. 2003) disclose that the dark matter is cold (CDM).

Direct detection experiments (Bertone, Hooper, & Silk 2005) impose strict upper constraints on the DM particle-nucleon scattering cross section. Therefore, the dark matter presence does not affect normal matter propagation through it; it only makes a contribution to the total mass of the Galaxy and consequently to the large-scale gravitational field.

However, this conclusion is valid only if the dark matter distribution is uniform. At the same time, contemporary cosmological conceptions suggest that the dark matter has very likely a clumpy structure. If the dark matter contains clumps, their small-scale gravitational field exerts extra influence on normal substance. For instance, if a stream of gas flows through clumpy dark matter, the field disturbs the flow. It results in momentum exchange between the stream and the dark matter and partial transformation of the kinetic energy of the stream into heat.

The dark matter perturbations played an important role in the universe structure formation (Gorbunov & Rubakov 2009). The clumps were formed in the early Universe from some primordial fluctuations, and their present-day mass distribution depends on the fluctuation spectrum that is not known very well. Inflation theory predicts that the spectrum had a flat Harrison-Zeldovich shape (Gorbunov & Rubakov 2009); so perturbations with all masses existed in the early Universe. However, the clump mass distribution may have a cut-off in the area of small masses (very small
clumps should be destroyed by free-streaming). The minimal possible clump mass depends on the physical nature of the dark matter particles (especially on the mass). For one of the most popular dark matter candidates — neutralino, the lightest SUSY particle — the limit estimations vary from $10^{-12} M_\odot$ (Zybin, Vysotsky, & Gurevich 1999) to $10^{-8} M_\odot$ (Hofmann, Schwarz, & Stocker 2001). However, clumps with masses near the limit should prevail, for the number of clumps grows with reduction in the mass.

Another important process strongly affecting the dark matter structure is the tidal destruction of small clumps by the bigger ones (Berezinsky, Dokuchaev, & Eroshenko 2006). As a consequence, a significant part of dark matter is not included in clumps forming a more or less uniform component. The presence of the uniform component in the dark matter distribution does not influence, of course, the small-scale structure of the gravitational field that is totally created only by density perturbations over the minimal level (i.e. by clumps). On the other hand, it is shown (Berezinsky, Dokuchaev, & Eroshenko 2006) that the densest central part of small clumps can hardly be destroyed by the tidal forces.

Hence if the dark matter is really composed of WIMPs, the presence of clumps of various masses down to very small is a necessity rather than a theoretical possibility. The biggest clumps can be directly observed: these are nothing else than galactic haloes. Smaller clumps are presently beyond experimental detection. Meanwhile, any evidence for their existence could be very important: for instance, if the low-mass clumps are found to be lacking, it will probably mean that the dark matter is not composed of WIMPs. The second reason while the presence of small clumps is important is that if the dark matter annihilation is possible, it goes mainly in clumps (Berezinsky, Dokuchaev, & Eroshenko 2006). In realistic cosmological models low-mass clumps collapse earlier and turn out to be significantly denser than the massive ones, and it is these clumps that give the main contribution to the annihilation signal. Thus, the space distribution of low-mass clumps is very important for indirect detection experiments.

As we will see, the drag effect acting on the clump always grows with the clump mass, and the distribution of the smallest clumps can hardly be disturbed by any real astrophysical gas stream. On the other hand, clumps form a sort of hierarchical structure, in which small objects are bound in the gravitational field of the large ones. As a result, small clumps can be dragged together with the heavy objects to which they belong.

### 2 THE STRUCTURE OF CLUMPY DARK MATTER

First of all, density and gravitational potential distributions of the clumpy matter need to be ascertained. Unfortunately, the problem remains to be solved. Numerous models of the dark matter structure and clump density profiles are considered in literature. In this article, we will use the results of (Bullock et al. 2001) as one of possible scenarios, escaping the discussion of which of the models of the structure formation is the most plausible, which is far beyond the scope of this work. Moreover, as we will see, the result is not very sensitive to the parameters of the density profile.

In accordance with (Bullock et al. 2001), we adopt the Navarro-Frenk-White density profile of a clump

$$\rho_{NFW} = \frac{\rho_s}{(r/r_s)(1 + r/r_s)^2}$$  \hspace{1cm} (1)

where $\rho_s$ and $r_s$ are characteristic ‘inner’ density and radius of the clump. It is worth mentioning that profile (1) cannot be valid for all $r$, as the total clump mass diverges when $r \to \infty$.

Instead of $\rho_s$ and $r_s$, it is convenient to use the virial radius $R_{vir}$, mass $M_{vir}$, and the concentration parameter $c_{vir}$

$$c_{vir} \equiv \frac{R_{vir}}{r_s}$$  \hspace{1cm} (2)

The above-mentioned quantities are related by the equations (see (Bullock et al. 2001) for details)

$$M_{vir} = \frac{4\pi}{3} \Delta_{vir} \rho_s R_{vir}^3$$  \hspace{1cm} (3)

$$M_{vir} = 4\pi \rho_s r_s^3 A(c_{vir})$$  \hspace{1cm} (4)

$$A(c_{vir}) \equiv \ln(c_{vir} + 1) - \frac{c_{vir}}{c_{vir} + 1}$$

Here $\rho_s$ is the mean universe density, $\Delta_{vir}$ is the virial overdensity. For the present epoch $\Delta_{vir} \approx 337$. From equations (4) we can see that among the quantities $\rho_s$, $r_s$, $R_{vir}$, $M_{vir}$, and $c_{vir}$ there are only two independent. It is convenient to use $M_{vir}$ and $c_{vir}$.

The clump potential is:

$$\phi(r) = \frac{G M_{vir}}{r_s A(c_{vir})} \ln \left(1 + \frac{r}{r_s} \right)$$  \hspace{1cm} (5)

It is worth noting that, contrary to the Newtonian case, the potential remains finite at the centre of the clump.

$$\phi_0 = \frac{G M_{vir}}{r_s A(c_{vir})}$$  \hspace{1cm} (6)

Equation (5) for the clump gravitational field cannot be valid for an arbitrary big radius $r$. First, as we have already mentioned, the total mass of a clump with profile (1) diverges when $r \to \infty$. Second, equation (5) is certainly inapplicable at a radius $r$ if $M < \frac{4\pi}{3} \rho_s r^3$, where $M$ is the clump mass and $\rho$ is the dark matter density averaged over a region vastly larger than the clump size (for instance, on the outskirts of the Solar System $\rho \simeq 0.3$ GeV/cm$^3$ (Gorbunov & Rubakov 2008)). Indeed, this condition means that the matter not included into the clump already prevails inside the sphere of radius $r$ circling the centre of the clump and gives a bigger contribution to the gravitational field at this radius. From the latter consideration, we can roughly estimate the radius $R_\Sigma$ of the “sphere of influence” of a clump, where its gravitational field of the clump dominates, and equation (5) is applicable $R_\Sigma > r_s$. Moreover, profile (1) obviously does not hold for $r > R_\Sigma$ because of the tidal perturbations, and the question of exact density distribution is extremely complex. Consequently, supposition $M \equiv M_{vir}$ is quite acceptable in our
approximative calculations. It is convenient to use
\[ \Xi \equiv \frac{R_\Sigma}{\tau_s} = \frac{3}{4\pi} \frac{M_{vir}}{\bar{v}^2 c_{vir}} = \frac{3}{2} \frac{\Delta_{vir} \rho_c}{\rho} \] (7)
instead of \( R_\Sigma \).

### 3 CLUMP INTERACTION WITH A STREAM OF GAS

The nature of the clump interaction with a stream of gas depends on the relation between \( R_\Sigma \) and the free length \( l_\Gamma \) of the gas atoms. We start our consideration from the case when \( R_\Sigma \ll l_\Gamma \). Then we can look upon the gas flow as being a stream of noninteracting particles and consider propagation of each atom separately. A dynamical friction between the clump and the stream appears as a result of the particle scattering in the gravitational field of the clump (Chandrasekhar 1943; Binney & Tremaine 2008).

Let us consider a stream of particles of mass \( m \) moving at a speed \( \bar{v}_\infty \) at infinity and scattered on potential \( \Xi \). The total momentum change of the gas stream in a unit time (i.e., the force \( \overline{\Xi} \) acting on the clump) is:
\[ \overline{\Xi} = \int_0^{R_\Sigma} \Delta p_s \cdot \bar{v}_\infty \cdot 2\pi \bar{w} d\bar{w} \] (8)
where \( \overline{\Xi} \) is the impact parameter of the trajectory and is \( n \) is the gas concentration. The problem can be simplified if we take into account that the clump gravitational field is respectively weak, and the scattering angle \( \theta \) is small. Then after simple calculations we obtain for the clump acceleration \( \bar{w} = \overline{\Xi} / c_{vir} \):
\[ w = \frac{\pi \nu G^2 M_{vir}}{3\nu^2 c_{vir}^2 A^2 (c_{vir})} \cdot \ln^3 \left( \frac{c_{vir}}{2} \frac{3}{2} \frac{\Delta_{vir} \rho_c}{\rho} \right) \] (9)
We introduced the gas density \( \nu = \bar{n}_\infty \). One can see that the result is not very sensitive to the precise determination of \( c_{vir} \) or \( \Xi \) while \( c_{vir} \gg 1 \) and \( \Xi \gg 1 \).

In order to complete the calculations, we should define the dependence \( c_{vir}(M_{vir}) \) in an explicit form. Unfortunately, there are no reliable estimations of this relationship in the range of small clump masses. (Bullock et al. 2001) reported about power-law growth of halo concentration \( c_{vir} \) with \( M_{vir} \) decreasing, and the results of (Ahn & Komatsu 2002) confirm this conclusion. However, if we try to interpolate this dependence to the area of the lowest possible concentration, we get grotesquely huge values of \( c_{vir} \). It is likely that in this realm the relationship is somehow modified; however, halo concentration should be very big (\( c_{vir} \gg 1 \)). Following (Ahn & Komatsu 2002), we adopt that \( c_{vir} \) is constant \( c_{vir} = 70 \) at low masses \( M_{vir} \ll 10^{10} M_\odot \) (though this premise looks unlikely). Considering the uncertainty of \( c_{vir} \) determination and the above-mentioned weak dependence of the result on \( c_{vir} \) and \( \Xi \) we can neglect the difference between \( R_\Sigma \) and \( R_{vir} \) (i.e., we adopt \( \bar{\rho} = \Delta_{vir} \rho_c \), which is quite natural for galactic haloes). Then \( c_{vir} = \Xi \), and it is easy to see that, if \( c_{vir} \gg 1 \), \[ \ln^3 \left( \frac{c_{vir}}{2} \frac{3}{2} \frac{\Delta_{vir} \rho_c}{\rho} \right) \cong \ln \left( \frac{c_{vir}}{2} \right), \] and we can simplify equation (9):
\[ w = \frac{\pi \nu G^2 M_{vir}}{3\nu^2 c_{vir}^2} \cdot \ln \left( \frac{c_{vir}}{2} \right) \] (10)

### Interaction of clumpy dark matter with gas

Formally, equation (10) rapidly grows when \( \bar{v}_\infty \to 0 \). However, if \( \bar{v}_\infty \) is small the scattering angle \( \theta \) becomes large, and (10) is no longer applicable. In order to estimate the maximum possible clump acceleration \( w_{max} \), we should perform a more detailed treatment.

In a general way, equation (8) can be rewritten as
\[ \overline{\Xi} = 2\nu \bar{v}_\infty^2 \int_0^{R_\Sigma} (1 - \cos \theta) \bar{w} d\bar{w} \] (11)
Since \( (1 - \cos \theta) \ll 2 \) we can affirm that
\[ \overline{\Xi} \approx 2\nu \bar{v}_\infty^2 R_\Sigma^2 \] (12)

On the other hand, when \( \bar{v}_\infty \) is very small, the scattering angle is large, and \( (1 - \cos \theta) \) is not a small number. Therefore, equation (12) can be considered as an estimation of \( \overline{\Xi} \) in the limit of small velocity. Consequently, the force grows as \( \nu^2 \) while \( \bar{v}_\infty \) is small and decreases as \( \nu^2 \nu \) for big \( \bar{v}_\infty \). We can estimate \( \nu^2 \) that provides the maximum of \( \overline{\Xi} \) equating expressions (9) and (12):
\[ \bar{v}_\infty^2 = \frac{GM_{vir}}{R_{vir}} \frac{\ln(c_{vir}/2)}{6} \] (13)
Substituting this equation into (12), we obtain
\[ w_{max} = 2\nu GR_{vir} \frac{\ln(c_{vir}/2)}{6} \] (14)
Maximum clump acceleration weakly (as \( \nu M_{vir} \)) depends on the clump mass, and if \( c_{vir} = 70 \), \( w_{max} \approx 4.8\nu GR_{vir} \).

In the case when \( R_\Sigma \gg l_\Gamma \), the hydrodynamic approach should be used. The streamline picture about the clump depends on the Mach number at infinity \( M_{vir} \equiv \bar{v}_\infty /a \), where \( a \) is the sound speed. It is extremely difficult to calculate the flow in potential field \( \Xi \) exactly. However, the problem is thoroughly studied numerically, and we can easily obtain simple estimations (Ostriker 1999; Sánchez-Salcedo & Brandenburg 1999). Beside this, we can use a similarity between the system under consideration and gas accretion on a compact astrophysical object: the later problem has been thoroughly studied (Lipunov 1992).

The case \( \nu_\infty < a \) corresponds to the spherical accretion (Bondi & Hoyle 1952), the case \( \nu_\infty > a \to 0 \) to the cylindrical one (Bondi & Hoyle 1944). We emphasize, however, the above-mentioned important distinction from the accretion problem: potential \( \Xi \) is everywhere finite \( |\phi| \leq |\phi_0| \).

If \( \nu_\infty \gg a \) we may neglect the pressure. Then the field of flow velocities coincides with the collisionless case, except for a narrow zone behind the clump, where the streamlines cross and a shock wave appears (see Bondi & Hoyle 1944 for details). The form of the shock is not quite clear, and we adopt the simplest supposition that it is a cone with the vertex at the centre of the clump and the corner angle \( \nu_\infty a = \nu_\infty \frac{1}{2} \bar{w}_{max} \) (Lipunov 1992), which corresponds to the solid angle \( d\Omega = \pi \left( \frac{\bar{w}}{\bar{w}_{max}} \right)^2 \). So all the space around the clump can be divided on two regions: region I before the shock wave, where the stream velocity coincides with the collisionless case, and conic region II after the shock wave. Consequently, the total force acting on the clump (and consequently, its acceleration) is a sum of its interactions with the flow in regions I and II (\( w_1 \) and \( w_2 \), respectively). In order to calculate \( w_1 \) we may use equation (10) as in the collisionless case.
To estimate $w_2$ we assume that the shock wave is strong. Then the gas after the shock wave is $\frac{1}{\gamma + 1}$ times denser than before, and the overdensity in region $II$ is $\nu_{\rm over} = \frac{1}{\gamma} \nu$. The total force of attraction $\vec{F}_2$ between the clump and the substance behind the wave can be roughly estimated as
\[
\vec{F}_2 = \int_0^{R_2} G M_{\rm vir} dm \frac{d\Omega}{r^2}
\]
(15)
where $dm = \nu_{\rm over} \sqrt{2} \frac{d\Omega}{r^2 \nu^2}$. Substituting here the equation for $d\Omega$, we obtain after simple calculations:
\[
w_2 = \frac{2\pi}{\gamma - 1} G \nu R_{\rm vir} \frac{a^2}{\nu_{\rm c}}
\]
(16)
It is interesting to compare the contributions $\vec{F}_1$ and $\vec{F}_2$. Dividing (10) by (16), we obtain:
\[
\frac{w_1}{w_2} = \frac{(\gamma - 1)}{6} \frac{A(c_{\rm vir}) \ln(c_{\rm vir}/2)}{c_{\rm vir}} \frac{\phi_0}{a^2}
\]
(17)
One can see that the ratio $w_1/w_2$ grows with $M_{\rm vir}$. The equality is reached at
\[
M_{\rm vir} = \frac{12a^3 \sqrt{2 \pi \sigma_{\nu} \rho_{\nu}}}{(G(\gamma - 1) \ln(c_{\rm vir}/2))^{3/2}}
\]
(18)
For instance, if $c_{\rm vir} \equiv 70$, $a = 1\,\text{km/s}$, $\gamma = 4/3$, it corresponds to $M_{\rm vir} \approx 3 \times 10^8 \, M_{\odot}$. So component $\vec{F}_2$ dominates for all the clumps except for the largest ones.

It is easy to see that in the opposite case, when $v_\infty \ll a$, the flow pattern is determined by the ratio between the minimum of the clump potential and $a^2$. If $|\phi_0| > a^2$, the drag force acting on the clump is negligible. In fact, in this case the flow is everywhere subsonic. Then, because of D’Alembert’s paradox, the drag force is totally created by the viscosity of the stream substance, that is always very small in astrophysical systems. If $|\phi_0| > a^2$ for $M_{\rm vir} > 2.5 \times 10^8 \, M_{\odot}$, if we adopt $c_{\rm vir} \equiv 70$, $a = 1\,\text{km/s}$. Consequently, small clumps do not interact with the stream.

If $|\phi_0|$ is big enough, the gas flow becomes supersonic at some region, and a shock wave may appear there. The wave occurrence leads to a strong enhancement of the drag force. By analogy with the previous case we can conclude that only the stream lines crossing the shock wave give a significant contribution into the drag force. The exact hydrodynamical solution of the task can hardly be found. However, we can easily estimate the force. It is reasonable to assume that the shock may appear only on the lines passing through the region where $|\phi| > a^2$. Its radius $r_c$ is determined by equation \(5\)
\[
\phi(r_c) = -\frac{G M_{\rm vir}}{r_c A(c_{\rm vir})} \frac{\ln(r_c/r_s + 1)}{r_c/r_s} = a^2
\]
(19)
A good approximation for $r_c$ is
\[
r_c = r_s \frac{\phi_0}{a^2} \ln \left( \frac{\phi_0}{a^2} \right)
\]
(20)
Now we may roughly estimate the force acting on the clump as $\frac{\nu v_c^2}{2} \pi r_c^2$
\[
\vec{F} = \frac{\pi}{2} \frac{\nu v_c^2}{2} \frac{\phi_0}{a^2} \ln \left( \frac{\phi_0}{a^2} \right)
\]
(21)
As we can see, the drag force increases as $v_\infty^2$ if $v_\infty \ll a$ and decreases as $v_\infty^{-2}$ if $v_\infty \gg a$. Consequently, we may assume that the force mounts to the maximum value when $v_\infty \approx a$.

Reviewing equations (10), (16) and (21), we can make another very common conclusion: the drag effect always grows with the clump mass. In fact, it is easy to see that the clump acceleration $w = \vec{F}/M_{\rm vir}$ increases with $M_{\rm vir}$ in each case considered. However, the shape of the growth is different. In the most important instance of a supersonic stream and large clumps the acceleration increase is relatively slow ($w \sim M_{\rm vir}^{1/3}$) if the clump mass is below the limit (18). For heavier clumps the acceleration is proportional to the mass ($w \sim M_{\rm vir}$), see Fig. 1.

**Figure 1.** The clump acceleration dependence on the clump mass. The force acting on the clump is determined by the sum of (10) and (16). We take $\nu = 10^{-25} \, \text{g/cm}^3$, $v_\infty = 1\,\text{km/s}$, $a = 10\,\text{km/s}$, $\gamma = 4/3$, $c_{\rm vir} = 70$.
\( \nu \simeq 2 \times 10^{-28} \text{g/cm}^3 \) \cite{Spitzer1978}. However, the drag force acting on the clump does not depend on the gas temperature: indeed, the interstellar gas is in hydrodynamical equilibrium, and its density \( \nu \propto T^{-1} \propto a^{-2} \). Consequently, the multiplier \( \nu a^2 \) in \ref{eq:10} remains constant. For definiteness sake we consider a hot cloud (\( a \simeq 10 \text{ km/s} \)). Taking the velocity difference between the halo and the disk rotations to be equal to \( v_\infty \simeq 150 \text{ km/s} \), we obtain \( w \simeq 8 \times 10^{-15} \text{ m/s}^2 \).

The age of the Galaxy is \( \sim 10^{10} \) years and even if the clump has never left the disk, its velocity change does not exceed \( \delta v \simeq 30 \text{ km/s} \).

Let us consider formation of our Galaxy \cite{BinneyMerrifield1998}. As of now, the details of this process are unclear. The Galaxy was likely formed as a result of the protogalaxy collapse, which size was at least an order larger than the present radius of the Galaxy disk. A merging of smaller structures played an important role in the process.

The collapse initially comes about almost as a free-falling. Subsequently, however, dissipation processes in the gas resulted in the separation of the normal and the dark matters. The collisionless dark matter stopped collapsing and formed an extensive and almost spherically symmetric halo. The gas component lost its energy via emission, kept on compressing and formed a compact thin disk. Consider if a significant momentum transition could take place during the collapse.

We set the characteristic radius where the dark matter detached to be equal to \( r \simeq 30 \text{ kpc} \simeq 10^{23} \text{ cm} \), the baryon mass of the Galaxy \(-\) to \( 3 \times 10^{11} M_\odot \simeq 6 \times 10^{44} \text{ g} \). Then the average density of the gas was equal to \( \nu \simeq 1.5 \times 10^{-25} \text{ g/cm}^3 \).

We adopt the protogalaxy temperature \( \sim 10^4 \text{ K} \), which corresponds to sound speed \( a \simeq 10 \text{ km/s} \). Substituting this value (and \( \gamma = 4/3 \), \( c_{\text{vir}} = 70 \)) to \ref{eq:13}, we obtain \( M_{\text{vir}} = 3 \times 10^6 M_\odot \). Below this mass we should use equation \ref{eq:10}. We adopt the time of separation of the baryon from the dark matter to be equal to \( t = 0.2 \text{ col} \), where \( t_{\text{col}} \) is the characteristic time scale of the collapse of the Galaxy \( (t_{\text{col}} \sim 10^8 \text{ years}) \). Substituting all the values into \ref{eq:10}, we finally obtain:

\[
\delta v = \frac{2\pi}{\gamma - 1} \frac{GmR_{\text{vir}}}{v_\infty^2} \simeq \ldots
\]

\[
\simeq 0.5 \left( \frac{a}{v_\infty} \right)^2 \left( \frac{M_{\text{vir}}}{M_\odot} \right) \frac{1}{10^3} \left[ \frac{\text{km}}{\text{s}} \right] \tag{22}
\]

Since the velocities of the dark and baryon matters were identical at the beginning of the collapse, and the drag force reaches its maximum when the velocity difference is of the order of the sound speed, we can estimate the minimal mass of the clump that could be dragged by the baryon matter by substituting \( v_\infty = a \) into \ref{eq:22}.

\[
\delta v_{\text{max}} \simeq 0.33 \left( \frac{\nu}{10^{-28} \text{ g/cm}^3} \right) \left( \frac{M_{\text{vir}}}{M_\odot} \right) \frac{1}{10^3} \left[ \frac{\text{km}}{\text{s}} \right] \tag{23}
\]

Taking \( \delta v = 50 \text{ km/s} \) as a significant velocity change, we obtain \( M_{\text{vir}} \sim 10^8 M_\odot \). A velocity change \( \delta v = 10 \text{ km/s} \) corresponds to \( M_{\text{vir}} \sim 10^4 M_\odot \). In all the cases the clump should be massive enough to be carried along by the gas. On the other hand, even much heavier clumps should be present during the Galaxy formation \( (M_{\text{vir}} \geq 10^7 M_\odot) \).

The momentum transmission from the gas to the dark matter leads to at least three important consequences. First, the large clumps, carried along by the gas, were destroyed by tidal forces. However, the dark matter of the mergers flowed into the halo, forming a more compact, oblate, and faster rotating substructure (usually called the thick disk). This effect has been already discovered numerically by \cite{Read2008}.

Second, the thick disk should be more clumpy than the halo. In fact, the gas collapse could carry along only the clumpy component of the halo, while the homogeneous remained almost spherically symmetric. Of course, the gas could drag only large clumps that were later destroyed by tidal forces. However, the large mergers contained a hierarchical system of smaller clumps, which were moved in such a way to the thick disk. These low-massive clumps are too small to be destroyed by tidal forces, and the thick disk turns out to be enriched with them. Meanwhile, these are the small clumps that give the main contribution to the possible dark matter annihilation signal, even if they make up only a small fraction of the dark matter. Thus we can expect that if some indirect dark matter search find any signal, the so-called boost factor \( C \equiv \langle \rho_{dm}^2 \rangle / \langle \rho_{dm} \rangle^2 \) will be higher in the thick disk.

Third, only the clumps situated relatively close to the Galaxy centre could be carried along with the gas. Indeed, the average density of the protogalaxy increased with its collapse as \( r^{-3} \). Consequently, the drag force also rapidly decreases with the radius. By analogy with equation \ref{eq:23} we can estimate the minimal mass of a clump, which velocity could be changed by the drag force on more than \( \delta v = 10 \text{ km/s} \) during the protogalaxy collapse. We determine the force acting on the clump as the sum of \ref{eq:10} and \ref{eq:16}, substituting there \( v_\infty = a \simeq 10 \text{ km/s} \) and the above-mentioned parameters of the protogalaxy. The results are represented in Fig. 2. We can see that above \( r = 150 \text{ kpc} \) the velocity disturbance is less than \( 10 \text{ km/s} \) even for the hugest clumps \( M_{\text{vir}} \geq 2 \times 10^8 M_\odot \). It means, that the dark matter structure at the outskirts of the halo remained un-

\( © 2011 \text{ RAS, MNRAS 000} \)
touched, and the big clumps from there were not transported to the thick disk. Consequently, the boost factor of the outer regions should be higher.

It is interesting to compare the described above effect with the efficiency of direct collisions of baryons with the dark matter particles. The number of collisions occurring in a unit volume in a unit time is \( \sigma n dm \frac{v}{\infty} \), each of them transmits a momentum \( \sim m_b v_{\infty} \) on the average. Here we have symbolized the concentration and mass of the dark matter particles and baryons by \( n_{dm}, m_{dm} \) and \( n_b, m_b \) respectively, \( \sigma \) is the cross-section of a dark matter particle scattering on a baryon. The mass of dark matter in the volume considered is \( m_{dm} n_{dm} \), and its acceleration is:

\[
w \sim \frac{\sigma n_{dm} v_{\infty} \cdot m_b v_{\infty}}{m_{dm} n_{dm}} = \sigma n_b v_{\infty}^2 \frac{m_b}{m_{dm}}
\]  

For a neutralino of mass 100 GeV a cross-section higher than \( 10^{-44} \) cm\(^2\) is now experimentally excluded \cite{Aprile2011}. The effect grows with the increasing of the relative speed of the colliding objects. However, even for the above-mentioned example of cluster collision, where the relative speed is the largest, acceleration \( \sim 2 \cdot 10^{-34} \) m/s\(^2\), which is at least eight orders lower than the acceleration produced by the mechanism under consideration.

5 ACKNOWLEDGEMENTS

This work was supported by the RFBR (Russian Foundation for Basic Research, Grant 08-02-00856).

REFERENCES


