ANGULAR MOMENTUM TRANSFER AND LACK OF FRAGMENTATION IN SELF-GRAVITATING ACCRETION FLOWS

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ABSTRACT

Rapid inflows associated with early galaxy formation lead to the accumulation of self-gravitating gas in the central regions of proto-galaxies. Such gas accumulations are prone to non-axisymmetric instabilities, as in the well-known Maclaurin sequence of rotating ellipsoids, which are accompanied by a catastrophic loss of angular momentum. Self-gravitating gas is also intuitively associated with star formation. However, recent numerical simulations of the infall process display highly turbulent flows, which do not fragment. We propose that the angular momentum transfer process, which enables the inflow, also suppresses fragmentation. We argue that inefficient angular momentum loss by the gas leads to decay of the turbulence, which in turn triggers global instabilities in the gas and renewed turbulence driving. We show that a flow regulated in this way is stable against fragmentation, whilst staying close to the instability threshold for bar formation — thick self-gravitating disks are prone to global instabilities before they become unstable locally. On smaller scales, the fraction of gravitationally unstable matter swept up by shocks in such a flow is a small and decreasing function of the Mach number. We therefore conclude, counterintuitively, that gas able to cool down to a small fraction of its virial temperature will not fragment as it collapses. This provides a venue for supermassive black holes to form via direct infall of gas into galactic nuclei, without the intermediary stage of forming a star cluster. Some black holes could have formed or grown extensively in massive dark matter halos at relatively low redshifts.

Subject headings: cosmology: dark matter — galaxies: evolution — galaxies: formation — galaxies: halos — galaxies: interactions — galaxies: kinematics and dynamics

1. INTRODUCTION

The formation of luminous galactic nuclei and supermassive black holes is inherently related to the gravitational collapse of gas and its concurrent release of large quantities of angular momentum. Recent numerical simulations of this process (Regan & Haeblnelt 2009; Levine et al. 2008; Wise, Turk & Abel 2008; Engenhaier & Shlosman 2004) reveal a number of puzzling empirical details. First, the collapse appears to be self-similar, over as many as 12 decades in radius, with a density profile of $\rho \propto r^{-2}$ when averaged over spheres, where $r$ is the spherical radius. This is observed despite an extremely asymmetric and chaotic appearance locally. Second, it is accompanied by fully developed, highly supersonic turbulence, with near-virial velocities at all radii. And third, it does not show any fragmentation, despite being self-gravitating and isothermal. The gas temperature is stable at a few×10$^3$ K (sound speed $c_s \sim 3 \text{ km s}^{-1}$), as governed by the atomic cooling rate, and the flow Mach number is $M \sim$ few. Here we propose an explanation for these phenomena.

Angular momentum transfer is of paramount importance to the evolution of many astrophysical systems. Such systems, which are kept intact by their internal gravitational energy, or self-gravity, are found on all scales, from planetary and stellar, to galactic and beyond. It is understood that (long-range) gravitational torques rather than (local) viscous torques are responsible for angular momentum redistribution in these systems (e.g., Lynden-Bell & Kalnajs 1972), especially on scales where torques generated by ordered magnetic fields (Blandford & Payne 1982) are inefficient.

The amplitude of gravitational torques in a gaseous self-gravitating object is determined by the gravitational quadrupole moment, i.e., by its departure from axial symmetry. When such an object is supported by rotation against gravitational collapse, its axial symmetry is known to be spontaneously broken when the ratio of bulk kinetic energy to (absolute) gravitational potential energy, $T/\left|W\right|$, is larger than some critical value. The Maclaurin sequence of rotating fluid ellipsoids, being an angular momentum, $J$, sequence, can serve as a simple quantitative example (see, e.g., Chandrasekhar 1969, for a review). At some value of $J$, the axial symmetry of an ellipsoid is broken abruptly by the $m = 2$ mode to a bar-like shape, due to the lower energy of this configuration. A more robust stability indicator, which is applicable to differentially rotating, centrally concentrated fluids, is $\alpha \equiv (T/\left|W\right|)/(\Omega/\Omega_1)$, where $0 \leq \alpha \leq 1/2$, $\Omega$ is the angular velocity and $\Omega_1 = (2\pi G \rho A )^{1/2}$ is the Jeans frequency, $A$ being the shape factor defined in terms of the meridional eccentricity (Christodoulou, Shlosman & Tohline 1995). The bar instability appears to be universal and is triggered in any self-gravitating system, collisional or collisionless, with various mass density profiles and sufficient bulk rotation. The subsequent evolution, however, depends on the ability of the system to dissipate energy.

When self-gravitating systems shed angular momentum, their potential (i.e., self-gravitating) and total energies can decrease as well. The relative amount of angu-
lar momentum loss, compared to potential energy loss, is crucial to determining the outcome. If this ratio is low enough, there is again a tendency to break axial symmetry and increase the quadrupole moment, as in the Maclaurin sequence. In non-dissipative systems, of course, the total energy is conserved. Such systems cannot be driven into a state of gravitational collapse by the loss of angular momentum, because the loss of potential energy saturates. On the other hand, in dissipative systems, energy loss can keep pace with the loss of angular momentum, forcing them into continuous collapse. The latter situation applies in fluid systems, such as gas-dominated disks, and forms the essence of the ‘bars-in-bars’ mechanism of angular momentum transfer in nested stellar/liquid or fluid/fluid bars, which is accompanied by dynamical inflow (Shlosman, Frank & Begelman 1989; Shlosman, Begelman & Frank 1990). Such fluid bars can form in maximally rotating neutron stars (e.g., Shibata, Baumgarte & Shapiro 2000), trigger the formation of black holes (Begelman, Volonteri & Rees 2006) or dominate galactic proto-disks at high and intermediate redshifts (e.g., Heller, Shlosman & Athanassoula 2007; Romano-Diaz et al. 2008), and hence provide a universal channel for angular momentum redistribution in self-gravitating systems.

2. TURBULENCE DRIVING

To ensure continuous collapse, a rotation-dominated self-gravitating system must efficiently decrease its $J$ and remain dissipative. Stellar systems (i.e., disk galaxies) retain the first property, but plainly lack the second. What is less clear is whether gasdynamical systems are guaranteed to remain dissipative as they collapse. Being dominated by self-gravity, they can fragment and these fragments can contract, sharply decreasing their interaction cross-sections and becoming essentially collisionless. Star formation is an example of this process. The internal dissipation will be dramatically curtailed by the star formation and, consequently, the overall collapse of the system will be terminated or at least delayed. However, the ability of fragments to contract depends on their equation of state and should not be taken for granted (Paczynski 1978; Shlosman et al. 1990). The requirement that inflow accompanies the transfer of $J$ in nested fluid bars critically depends on the non-fragmentation of the flow. It has never been shown that the gasdynamical bars-in-bars mechanism can overcome this limitation.

The first hints that fragmentation may be avoidable in bars-in-bars collapse have been provided recently by numerical simulations, which reveal that nested bar instabilities accompany the gravitational collapse, as predicted. The flow encompasses two different dynamical regimes: (1) a dark matter (DM)-dominated regime, from the growing halo virial radius down to a few pc; and (2) a gas-dominated regime, within the central few pc. The latter regime is characterized by the dynamical decoupling of self-gravitating gas from the DM background (Englaier & Shlosman 2004). The characteristic radius of a few pc applies to DM halos with virial temperatures somewhat in excess of $10^4$ K. More massive halos will have correspondingly larger gas-dominated central regions, and hence the final product of gravitational collapse can be much more massive.

It is not surprising that supersonically inflowing gas should become turbulent but somewhat surprising that the turbulence should maintain a virial level throughout the flow, given its high dissipation rate — this requires an efficient driving mechanism. On larger scales, the turbulence can be driven by a variety of mechanisms inherent in the hierarchical buildup of structure in the Universe, such as mergers. On smaller scales, the turbulence can be driven by the gravitational energy released during the infall. As gas with some seed $J$ goes into gravitational collapse, the centrifugal barrier will slow it down at some cylindrical radius $R$. This will happen concurrently with the decay of turbulence and will drive $T/|W|$ over the instability threshold. The instability develops on a dynamical timescale. The growth of the barlike mode in the gas will excite shocks and force the gas to lose rotational support, forcing it into further gravitational collapse and pumping energy into turbulent motions. Thus, we suggest that the persistence of virial, supersonic turbulence in a self-gravitating inflow is the consequence of a self-regulating instability. This cascade will be quenched only when the gas becomes optically thick, resulting in the sharp rise of thermal energy compared to turbulent energy. We quantify the above processes in §3.

As an example, we assume a time-independent, self-similar, self-gravitating, isothermal disklike configuration with Keplerian velocity $v_K$ independent of radius. (The latter is a natural outcome for a non-rotating flow [Larson 1969], but it can be extended to inflows with substantial rotation.) The surface density scales as $R^{-1}$. If the inflow speed is roughly independent of radius (as suggested by the simulations), it does not affect the radial momentum equation (since $v_R \ddt R / dR \approx 0$). One can then construct a crude model for the radial structure of the inflow. Neglecting thermal pressure, one finds a tradeoff between the typical turbulent velocity $v_t$ and the rotation speed $v_\phi$, $v_\phi^2 = v_K^2 - 2v_t^2$. If $v_t$ is small, the rotation is close to Keplerian and the ratio $T/|W|$, which is roughly proportional to $(v_\phi / v_K)^2$, exceeds the threshold for bar instability. As the radial inflow slows down, the turbulence decays — to the point where $T/|W|$ approaches the stability threshold. The decreasing ratio $v_t / v_\phi$ will lead to self-similar vertical structure as well, with $\rho \propto R^{-2}$, as seen in the simulations. (We note that the self-similar, self-gravitating disk models of Toomre [1982] provide a promising platform for refining this picture.) This example also holds, with straightforward quantitative modifications, when the gas inflow proceeds in the background potential of the DM halo.

3. DAMPING FRAGMENTATION IN THICK DISKS

The next aspect of the flow that needs to be addressed is fragmentation. If $v_t \ll v_\phi$ and the thermal energy dominates over the turbulent one, one can show that the disk is strongly unstable to breakup into rings (Toomre 1964). A direct application of the Toomre criterion, which is a generalized Jeans instability applied to a razor-thin, differentially-rotating disks, shows that the disk becomes more stable as $v_t$ increases at the expense of $v_\phi$. We test whether a self-gravitating gaseous disk which is subject to the (global) bar instability is stable locally, i.e., whether it is subject to fragmentation. Here, we estimate the effect of the finite thickness induced by the turbulence on the local stability of the disk. In the following, we refer to 2-D disks as ‘thin’ disks, and to disks with a finite
thickness as ‘thick’ disks. The latter are fully 3-D rotating configurations which are supported vertically (along the rotation axis) against gravity by turbulent motions, while in the equatorial plane, both rotation and turbulent motions contribute to this support.

The dispersion relation for a razor-thin, 2-D gaseous disk with a surface density \( \Sigma \) and a turbulent velocity \( v_t \) (which replaces the sound speed) is given by

\[
\omega^2 = \kappa^2 + v_t^2 k^2 - 2\pi G \Sigma |k|, \tag{1}
\]

where \( \kappa \) is the epicyclic frequency, \( k \) is the radial wavevector and \( \omega(k) \) is the wave frequency in the inertial frame. The isothermal disks considered here have a flat rotation curve and consequently \( \kappa^2 = 2\Omega^2 \), where \( \Omega \) is the angular velocity in the disk. In an unperturbed self-gravitating disk, the local Keplerian velocity in the disk is given by \( v_K^2 = v_o^2 + 2v_t^2 \), where \( v_o = \Omega R \). The tradeoff between the turbulent “pressure” and the rotational support against gravity, for a fixed \( v_K \), is crucial for understanding the disk dynamics and its stability.

Thin gaseous disks are stable to fragmentation if the Toomre parameter \( Q \) satisfies

\[
Q = \frac{\kappa v_t}{\pi G \Sigma} > 1. \tag{2}
\]

These disks are stabilized by rotation on larger scales and by the thermal (or turbulent) pressure on smaller scales. Sufficiently large \( Q \) ensures that these two scales overlap and the disk does not fragment. Both thermal and turbulent pressures in these disks are assumed to act in the equatorial plane only.

In reality, a disk with finite, isotropic pressure will have a finite vertical thickness, \( h \), which is maintained by the turbulence: \( h = (v_t/v_K)R \). The finite disk thickness affects the dispersion relation (eq. 1) by decreasing the value of the destabilizing gravity term (which appears with the negative sign), as follows:

\[
\omega^2 = \kappa^2 + v_t^2 k^2 - \frac{2\pi G \Sigma |k|}{1 + |k|h}. \tag{3}
\]

The correction factor, \( (1 + |k|h)^{-1} \), has been obtained by solving the Poisson equation along the rotation (z) axis, assuming that the disk has an exponential density profile, \( \Sigma(z) \sim e^{-|z|/h} \). Note that eq. (3) reduces to the usual dispersion relation for Jeans instability in the limit \( h/R \ll 1 \). The mass enclosed within a radius \( R \) by an isothermal, self-gravitating disk is given by \( M = 2\pi R^2 \), so that by using \( GM/R^2 = v_K^2 / R \), the dispersion relation (eq. 3) can be re-written as

\[
\frac{\omega^2 R^2}{v_K^2} = \frac{v_o^2}{v_K^2} + \frac{v_t^2}{v_K^2} (kR)^2 - \frac{kR}{1 + (kR)v_o/v_K}. \tag{4}
\]

Introducing dimensionless variables, \( y = kR \), \( x = (v_o/v_K)^2 \), and \( \zeta = \omega R/v_K \), so that \( (v_o/v_K)^2 = (1 - x)/2 \), we get

\[
\zeta^2 = 2x + \frac{1 - x}{2} y^2 - \frac{y}{1 + y(1 - x)}/2, \tag{5}
\]

which is the dispersion relation for self-gravitating, isothermal gaseous disks with a finite thickness contributed by turbulent motions. Our goal is to obtain the solution of this equation at the threshold for bar instability in such disks. The latter is given by Christodoulou et al. (1995) as

\[
\alpha \equiv \left( \frac{f T}{2 |W|} \right)^{1/2} > \alpha_{\text{crit}} \approx 0.35, \tag{6}
\]

where \( T = (v_o/v_K)^2 (v_K^4 R/2G) = x(v_K^4 R/2G) \) is the kinetic energy of the disk out to radius \( R \), \( |W| = f \sqrt{|GMdM/R|} = v_K^4 R/G \) is the potential energy contained within the same radius, and \( f \) is the shape factor. In our dimensionless units, \( T/|W| = x/2 \). The shape factor is \( f = 1 \) for a razor-thin disk and \( f = (2/3)(1 + T/|W|) \) for thick, differentially rotating spheroidal mass configurations. At marginal global stability of the disk, \( \alpha = \alpha_{\text{crit}} \); by substituting \( \alpha_{\text{crit}} \) into eq. (6), we obtain \( x_{\text{crit}} \approx 0.57 \).

Fig. 1 displays the local and global stability regimes in the \( x-y \) plane, i.e., \( (v_o/v_K)^2 - kR \), plane under imposed conditions as described below. The dark-grey regions delineate regimes of local instability, i.e., which are subject to fragmentation. In each frame, gas which lies in the left part of the diagram has negligible rotation, which
corresponds to the usual Jeans instability. The right part of each frame shows rotating turbulent gas, corresponding rather to Toomre’s instability. The left frame depicts stability conditions under the artificial assumption that the disk is 2-D (razor-thin). As expected, the right hand-side, for \( x > 0.49 \), is unstable over a wide range of wavelengths, because the turbulent pressure and rotation are not sufficient to prevent fragmentation. The second frame corresponds to a case in which we have (artificially) allowed the disk to maintain 5% of its thickness compared to what it should actually have based on \( v_t \). The unstable region has moved to \( x > 0.66 \). When the full value of disk thickness is accounted for, (right frame), the stability boundary lies at \( x \geq 0.94 \), i.e., in the regime of rapid rotation.

The light gray area shows the region of global (i.e., bar) instability calculated above, \( x > x_{\text{crit}} \approx 0.57 \). We observe that gas flows which possess fully developed turbulence must move from left to right in Fig. 1, as the turbulence decays with time. Hence the gas will first reach the threshold for bar instability while it is still locally stable. When the finite thickness of the disk is accounted for in the dispersion relation, the global (i.e., bar) instability sets in well before the disk is able to fragment.

What about fragmentation on small (sub-Jeans) scales, due to the compression of gas in the isothermal shocks that characterize supersonic turbulence? To address this, we adopt a widely used model of supersonic turbulence (e.g., Padoan 1995; Krumholz & McKee 2005; for reviews: Mac Low & Klessen 2004; Elmegreen & Scalo 2004; McKee & Ostriker 2007). We estimate the fraction of the material swept up by the shocks that becomes Jeans unstable, by following the formalism of Padoan & Nordlund (2002) and Krumholtz & McKee (2005). The turbulence is characterized by a lognormal probability distribution,

\[
p(x) = \frac{1}{2\pi \sigma_p^2} \frac{1}{x} \exp \left[ -\frac{(\ln x - \ln x_0)^2}{2\sigma_p^2} \right],
\]

with the mean \( \ln x = -0.5 \sigma_p^2 \), where \( x \equiv \rho / \rho_0 \) and \( \rho_0 \) is the mean density, and the dispersion \( \sigma_p \approx \ln(1 + 3M^2/4)\sigma_p^2 \).

Two spatial scales characterize such turbulence: the Jeans scale, \( \lambda_J = (\pi c_s^2/G\rho)^{0.5} \), and the transition scale, \( \lambda_s \), below which the flow becomes subsonic. The latter scale relates to the supersonic velocity dispersion, \( \sigma_t \), by \( \sigma_t \sim c_s (R/\lambda_J) \), which is valid over a large dynamic range: with \( \beta = 0.4 - 0.5 \) (Larson 1981; Krumholz & McKee 2005). In the overdense regions \( \lambda_J \sim \lambda_j \simeq 10^{-0.5} \lambda_J \), where \( \lambda_j \) is the Jeans length at \( \rho_0 \). To estimate the fraction of swept up gas that is Jeans-unstable \emph{over one free-fall time}, \( f_1 \), we integrate the lognormal distribution from \( x_{\text{crit}} \approx \rho_{\text{crit}} / \rho_0 = (\lambda_j / \lambda_s)^2 \approx 6 \gamma_{0.5,5} M^2 \) to infinity, where we have assumed \( \beta = 0.5 \) and \( \gamma_{0.5,5} \equiv 0.5v_t / v_K \approx 1 \). Strikingly, the dependence of \( f_1 \) on \( M \) assures that \( f_1 \lesssim 2 \times 10^{-2} \) for \( M \gtrsim 3 \) — such Mach numbers are routinely observed in simulations of inflows associated with galaxy formation. This differs from the gas evolution in galactic disks, where supersonic turbulence develops as well, but the turbulent velocities are very sub-\( \gamma \)-virial (because of lack of sufficient driving). Further high-resolution numerical simulations are needed to shed light on the intricacies of shock dissipation in supersonic turbulence over a wider range of Mach numbers.

4. DISCUSSION

To summarize, we have argued that the supersonic collapse of self-gravitating gas driven by nested gaseous bars is, counter-intuitively, not susceptible to fragmentation. Fragmentation is suppressed whenever the gas temperature remains below the virial temperature; moreover, the resistance to fragmentation increases with the infall Mach number. If the gas has sufficient time to achieve a quasi-equilibrium, when it has cooled down or its turbulence has decayed, it will be subject to fragmentation. If, however, a decrease in the level of turbulence leads to the renewed non-axisymmetric instability as we have argued, then the infall will recommence. The slowdown of radial inflow in regions of relative stability leads to spikes at which \( \alpha \) substantially exceeds \( \alpha_{\text{crit}} \), leading to the creation of a distinct bar and a rapid acceleration of infall. Thus, the bars-in-bars mechanism seems to be a highly intermittent one, as originally envisaged (Shlosman et al. 1989, 1990).

If confirmed by carefully designed numerical experiments, this scenario could have important implications for the formation and evolution of the central regions of galaxies and, in particular, the formation of supermassive black holes (SBHs). Begelman et al. (2006) suggested that the formation of SBH seeds by direct collapse would be suppressed in \( 10^4 \) K haloes that were metal-enriched or had significant cooling by H\(_2\), because the infalling gas would fragment and form stars. If this is not true, then a much larger fraction of \( 10^5 \) K halos could form single massive objects. We note that, within this framework, DM halos with larger virial temperatures could form more massive SBH seeds via direct collapse, if a sufficient amount of gas accumulates in their midst. Moreover, this mechanism would open the door to the formation or rapid growth of SBHs at later epochs in such massive halos, for example, following gas-rich mergers. There are also implications for the triggering of starbursts and the fueling of AGN in galactic nuclei where an SBH already exists.

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