How Do Disks Survive Mergers?

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ABSTRACT

We develop a general physical model for how galactic disks survive and/or are destroyed in mergers and interactions. Based on simple dynamical arguments, we show that gas primarily loses angular momentum to internal torques in a merger, induced by the gravity of the secondary. Gas within some characteristic radius, determined by the efficiency of this angular momentum loss (itself a function of the orbital parameters, mass ratio, and gas fraction of the merging galaxies), will quickly lose angular momentum to the stars sharing the perturbed host disk, fall to the center, and be consumed in a starburst. We use a similar analysis to determine where violent relaxation of the premerger stellar disks is efficient on final coalescence. Our model describes both the dissipative and dissipationless components of the merger, and allows us to predict, for a given arbitrary encounter, the stellar and gas content of the material that will survive (without significant angular momentum loss or violent relaxation) to re-form a disk in the merger remnant, versus being dissipationlessly and violently relaxed or dissipationally losing angular momentum and forming a compact central starburst. We test these predictions with a large library of hydrodynamic merger simulations, and show that they agree well (with small scatter) with the properties of simulated merger remnants as a function of merger mass ratio, orbital parameters, and gas distributions in simulations which span a wide range of parameter space in these properties as well as prescriptions for gas physics, stellar and active galactic nucleus (AGN) feedback, halo and initial disk structural properties, redshift, and galaxy masses. We show that, in an immediate (short-term) sense, the amount of stellar or gaseous disk that survives or re-forms following a given interaction can be understood purely in terms of simple, well understood gravitational physics, independent of the details of the interstellar medium gas physics or stellar and AGN feedback. This allows us to demonstrate and quantify how these physics are in fact important, in an indirect sense, to enable disks to survive mergers by lowering star formation efficiencies in low-mass systems (allowing them to retain large gas fractions) and distributing the gas to large radii. The efficiency of disk destruction in mergers is a strong function of gas content—our model allows us to explicitly predict and demonstrate how, in sufficiently gas-rich mergers (with quite general orbital parameters), even 1:1 mass-ratio mergers can yield disk-dominated remnants, and more realistic 1:3–1:4 mass-ratio major mergers can yield systems with < 20% of their mass in bulges. We discuss a number of implications of this modeling for the abundance and morphology of bulges as a function of mass and redshift, and provide simple prescriptions for the implementation of our results in analytic or semianalytic models of galaxy formation.

Key words: galaxies: active – galaxies: evolution – cosmology: theory

Online-only material: color figures

1. INTRODUCTION

In the now established “concordance” A cold dark matter (CDM) cosmology, structure grows hierarchically (e.g., White & Rees 1978), making mergers and interactions between galaxies an essential and inescapable process in galaxy formation. Indeed, mergers are widely believed to be responsible for the morphologies of spheroids (bulges in disks and elliptical galaxies; Toomre 1977), and observations find recent merger remnants in considerable abundance in the local universe (Schweizer 1982; Lake & Dressler 1986; Doyon et al. 1994; Shier & Fischer 1998; James et al. 1999; Genzel et al. 2001; Tacconi et al. 2002; Dasyra et al. 2006, 2007; Rothberg & Joseph 2004, 2006) as well as, e.g., faint shells and tidal features common around apparently “normal” galaxies (Malin & Carter 1980, 1983; Schweizer 1980, 1996; Schweizer & Seitzer 1992), which are thought to be signatures of galaxy collisions (e.g., Hernquist & Quinn 1988; Hernquist & Spergel 1992).

From both theoretical grounds (Ostriker & Tremaine 1975; Maller et al. 2006; Fakhouri & Ma 2008; Stewart et al. 2008, and references therein) and observations (e.g., Lin et al. 2004; Barton et al. 2007; Woods et al. 2006; Woods & Geller 2007) it appears that “minor” mergers of mass ratios $\lesssim 1:10$ are ubiquitous (there are almost no galaxies without mergers of at least this mass ratio in the last few Gyr), and moreover a large fraction ($\sim 1/2$ of the $\sim L_\star$ galaxy population is observed and expected to have experienced a “major” merger (mass ratio $\lesssim 1:3$) since $z \sim 2–3$ (Lotz et al. 2008; Bell et al. 2006; Bridge et al. 2007; Lin et al. 2008; Kartaltepe et al. 2007). With increasing redshift, kinematic and morphological indications of recent, violent disturbance in disk-dominated galaxies appear more frequent (Hammer et al. 2005; Flores et al. 2006; Puech et al. 2007, 2008).

Far from there not being enough mergers to explain the abundance of bulges and ellipticals, this has led to the concern that there may be far too many mergers to explain the survival and abundance of galactic disks in the context of our present understanding of galaxy formation. Toomre & Toomre (1972) were among the first to point out that mergers are capable of dramatically altering the morphologies of disks, transforming them into elliptical galaxies. Although their neglect of the importance of dissipational star formation and gas dynamics in the mergers led to some controversy (e.g., Ostriker 1980; Carlberg 1986; Gunn...
1987; Kormendy 1989), it is now increasingly well established that major mergers between spiral galaxies (similar to those observed locally and at $z \lesssim 2–3$) with gas fractions comparable to those observed yield remnants in good agreement with essentially all observed properties of low and intermediate-mass local elliptical galaxies (e.g., morphologies, shapes, sizes, kinematics, densities, colors, black hole (BH) properties, fundamental scaling relations, stellar populations, and halo gas; Hernquist 1989; Barnes & Hernquist 1991, 1996; Hernquist et al. 1993; Mihos & Hernquist 1994b, 1996; Di Matteo et al. 2005; Naab et al. 2006; Jesseit et al. 2007; Cox et al. 2006a, 2006b; Robertson et al. 2006b; Springel et al. 2005a; Burkert et al. 2008; Hopkins et al. 2007c, 2008a, 2008b, 2008c, 2008d).

Many intermediate and low-luminosity “cusp” ellipticals (encompassing ~80%–90% of the mass density in ellipticals) contain significant embedded disks (perhaps all such ellipticals, given projection effects; see Ferrarese et al. 1994; Lauer et al. 2005), and they form a continuous sequence with most SO galaxies, known to have prominent stellar (and even gaseous) disks (Kormendy 1983; Bender et al. 1992; Ferrarese et al. 1994, 2006; Kormendy et al. 1994; Lauer et al. 1995; Faber et al. 1997; Kormendy 1999; Emsellem et al. 2007). Indeed, the existence of embedded disks in simulated merger remnants is critical to matching the properties described above.

A wide variety of observations including stellar populations and star formation histories (e.g., Bender et al. 1989; Trager et al. 2000; McDermid et al. 2006) and kinematic and structural analysis of recent merger remnants (Schweizer et al. 1983; Schweizer 1983; Schweizer & Seitzer 1992, 2007; Hibbard & Yun 1999; Rothberg & Joseph 2004) demonstrate that most of these disks are not accreted in the standard cosmological fashion after the spheroid forms—they must somehow survive the merger or form very quickly thereafter from gas already in and around the galaxies. Therefore, despite the destruction of a large portion of a stellar disk in major mergers, some disk must survive mergers, and the amount that does so is a critical component determining many of the photometric and kinematic properties of even bulge-dominated and elliptical galaxies.

Moreover, “minor” mergers—at least those with mass ratios $\lesssim 10:1$ (below which the difference between “merger” and accretion becomes increasingly blurred)—are not generally believed to entirely destroy disks, but they are almost an order of magnitude more frequent than major mergers and as such may pose a more severe a problem for disk survival. In the $\Lambda$CDM cosmology, and from observed satellite fractions, it is unlikely that any disk (let alone a large fraction of disk galaxies) with a significant stellar age has survived $\sim 5–10$ Gyr without experiencing a merger of mass ratio $10:1$ or larger. Simulations (Quinn 1984; Quinn & Goodman 1986; Quinn et al. 1993; Hernquist & Mihos 1995; Walker et al. 1996; Velazquez & White 1999; Naab & Burkert 2003; Bournaud et al. 2005; Younger et al. 2007, 2008) and analytic arguments (Ostriker & Tremaine 1975; Toth & Ostriker 1992; Sellwood et al. 1998) suggest that gas-poor minor mergers can convert a considerable fraction of a stellar disk into bulge and cause significant perturbation (“puffing up” via dynamical heating) to the disk. The observed coldness of galactic disks suggests that this may be a severe problem; Toth & Ostriker (1992) argued that large disks such as that in the Milky Way could not have undergone a merger of mass ratio $\lesssim 10:1$ in the last $\sim 10$ Gyr. More recently, e.g., Stewart et al. (2008) and Hammer et al. (2007) emphasized that the tension between these constraints and the expectation in CDM models that a number of such mergers should occur imply either a deficit in our understanding of hierarchical disk formation or a challenge to the concordance cosmological model.

Given the successes of the $\Lambda$CDM model on large scales, and the increasing observational confirmation that disks do undergo (and therefore must somehow survive) a large number of mergers, it is likely that the problem lies in our (still relatively poor) understanding of disk galaxy formation. This has led to a great deal of focus on the problem of forming realistic disks in a cosmological context, with many different attempts and debate on the missing elements necessary to produce disks in simulations. Various groups have argued that self-consistent treatment of gas physics and star formation along with implementation of feedback of different kinds is necessary, along with greatly improved numerical resolution (Weil et al. 1998; Sommer-Larsen et al. 1999, 2003; Thacker & Couchman 2000, 2001; Abadi et al. 2003; Governato et al. 2004, 2007; Robertson et al. 2004; Okamoto et al. 2005; Scannapieco et al. 2008), in order to enable disks to survive their expected violent merger histories without completely losing angular momentum and transforming into systems that are too compact and have too much bulge mass (relative to real observed disks) by $z = 0$.

It has been known for some time (see Hernquist & Barnes 1991; Barnes & Hernquist 1996) that (even without any feedback) some fraction of the gas in even a major merger of two disks can survive and form new, embedded disks in the remnant, i.e., despite the problems outlined above, disks are not necessarily completely destroyed in mergers. However, early studies of this were restricted to cases with low gas content ($f_{\text{gas}} \lesssim 10\%$ in the progenitor disks), most of which was rapidly consumed in star formation, yielding small remnant disks in strongly bulge-dominated remnants. In seminal work, Springel & Hernquist (2005) and Robertson et al. (2006a) showed that, in idealized merger simulations with significant stellar feedback to allow the stable evolution of extremely gas-rich disks ($f_{\text{gas}} \sim 1$), even a major merger can produce a disk-dominated remnant. This has since been confirmed in fully cosmological simulations (Governato et al. 2007). Together with other recent investigations (see references above), these works have led to the growing consensus that a combination of strong stellar feedback and large gas content is essential for the survival of disk galaxies.

A large number of open questions remain, however. How, exactly, does feedback allow disks to survive mergers? What are the most important physics? Does it require fine-tuning of feedback prescriptions? How might things vary as a function of galaxy mass, redshift, gas content, merger orbits, and environment? Fundamentally, should this be expected for typical cosmological circumstances, or are these cases pathological?

The ambiguity largely owes to the fact that there is no deep physical understanding of how disks survive or re-form after mergers and interactions. It has only just become possible to conduct simulations with the requisite large gas fractions, and thus far theoretical explanations have largely been restricted to phenomenological analysis, with continued efforts to improve resolution and subresolution prescriptions. Moreover, without a full model for how disks behave in interactions, these simulations cannot be placed into the broader context of the emergence of the entire Hubble sequence (for example, asking the question, are the disks in lenticulars and embedded disks in ellipticals survivors of their premerger disks? Are they reaccreted? What determines how large they are? What is the key physics that gives rise to realistic embedded disks, leading to bulge-dominated galaxies with kinematic and photometric properties
similar to those in the real universe?) or within a fully cosmological context.

The resolution requirements for full models of disk formation are severe—limiting any attempt to properly simulate a cosmological box and still achieve the resolution necessary to reliably model a disk population—and so models of the population of disks, largely semianalytic, are forced to adopt simplified and untested prescriptions for the behavior of disks in mergers. This, in turn, has led to other well-known problems in modeling disk populations (even where prescriptions can ensure no artificial angular momentum losses); even when the cumulative (morphology-independent) galaxy mass function is correctly artificial angular momentum losses); even when the cumulative (morphology-independent) galaxy mass function is correctly predicted at the low-mass end, semianalytic models widely overpredict the relative abundance of low-mass spheroids and underproduce disks (even when satellites, which have other associated model uncertainties, are removed from consideration; see Somerville et al. 2001; Somerville et al. 2008; Croton et al. 2006; Bower et al. 2006; de Lucia & Blaizot 2007). Lacking a proper, physically motivated understanding of how low-mass or gas-rich disks may or may not survive mergers, attempts to address this problem in the models have been purely phenomenological and involve arbitrary prescriptions (see Koda et al. 2007).

Motivated by these concerns, in this paper we develop a physical, dynamical model for how disks survive and are destroyed in mergers and interactions. We show that, in an immediate (short-term) sense, the amount of stellar or gaseous disk that survives or re-forms following a given interaction can be understood purely in terms of simple, well-understood gravitational physics. Knowing these physics, we develop an analytic model that allows us to accurately predict how much of a given premerger stellar and cold gas disk will survive a merger, as a function of the merger mass ratio, orbital parameters, premerger cold gas fraction, and mass distribution of the gas and stars. We compare these predictions to the results of a large library of hundreds of hydrodynamic simulations of galaxy mergers and interactions, spanning a wide parameter space in these properties as well as prescriptions for gas physics, stellar and active galactic nucleus (AGN) feedback, halo and initial disk structural properties, redshift, and absolute galaxy masses. Our numerical experiments confirm that the analytic scalings accurately describe the behavior and bulge formation/disk destruction in mergers over the entire dynamic range surveyed, and confirm that the parameters not explicitly included in our model do not systematically affect either the mean predictions or the scatter of simulations about those predictions. This allows us to understand the mean behavior of systems with different orbits and mass ratios, as well as why systems with large gas fractions can form little bulge in even major mergers.

This is possible because gas, in mergers, primarily loses angular momentum to internal gravitational torques (from the stars in the same disk) owing to asymmetries in the galaxy induced by the merger. Hydrodynamic torques and the direct torquing of the secondary are second-order effects, and very inefficient. Once gas is drained of angular momentum, there is little alternative but for it to fall to the center of the galaxy and form stars, regardless of the details of the prescriptions for star formation and feedback (these may change things at the ~10%–20% level by blowing out some of the gas, but they cannot fundamentally alter the fact that cold gas with no angular momentum will be largely unable to form any sort of disk, or the fact that a galaxy’s worth of gas compressed to high densities and small radii will inevitably form a large mass in stars). But if the systems are sufficiently gas-rich, then there is little stellar material sharing the disk to torque on the gas in the interaction, and little or no angular momentum is lost.

Feedback can dramatically alter the ability of a disk to survive in a cosmological sense: by allowing galaxies to retain large gas fractions (as opposed to no-feedback scenarios, in which cold gas in a disk is usually quickly converted into stars), they are more gas-rich when they undergo interactions, allowing them to avoid angular momentum loss for the above reason. Moreover, we show that in detail (owing to the resonant structure of interactions), it is really gas within a certain radius of the stellar disk that is drained of angular momentum. The commonly invoked stellar wind feedback then enables cosmological disk survival in a second fashion: by redistributing gas out to large radii, it prevents angular momentum loss and allows rapid reformation of disks after a merger. Independent of any tuning, our model allows us to quantify the disks expected as a function of interactions of arbitrary properties, and to physically, explicitly quantify what the requirements are for feedback, in a cosmological scenario, to enable disk survival.

In Section 2, we describe our library of gas-rich merger simulations, which we use to test our physical model for disk destruction and survival. In Section 3, we demonstrate the existence of genuine disks in remnants of even major mergers and briefly consider their properties, and compare methods to separate the disks and bulges in merger remnants. In Section 4, we consider the question of how these disks form in and survive mergers: we identify the key components of any merger remnant in Section 4.1, highlighting that these disks originate from a combination of undestroyed premerger stellar disks and gas which avoids angular momentum loss in the merger. In Section 4.2, we discuss how, in detail, that angular momentum loss proceeds. We use this, in Section 4.3, to build a physical model for how angular momentum loss proceeds in mergers and predict the surviving disk content of merger remnants: we model and test how this depends on the gas content of the premerger disks (Section 4.3.1), and the orbital parameters (Section 4.3.2) and mass ratio (Section 4.3.3) of the encounter. We generalize to first passage and BH encounters (Section 4.3.4) and demonstrate that (for otherwise fixed conditions at the time of an encounter) our conclusions are purely dynamical, independent of feedback physics or details in our treatment of, e.g., star formation and the interstellar medium (ISM) gas physics (Section 4.3.5), although we use our model to determine exactly how these choices can have dramatic indirect consequences for disk survival (by altering the state of systems leading into a merger). We discuss some exceptions and pathological cases in Section 4.3.6, and relate our results to the long-term secular evolution of barred systems in Section 4.3.7. In Section 5, we outline how these results can and should be applied in analytic and semianalytic models of galaxy formation, and give appropriate prescriptions derived from our numerical experiments. Finally, we summarize our results and discuss some of their cosmological implications and applications to other models and observations in Section 6.

Throughout, we assume a Ω_M = 0.3, Ω_Λ = 0.7, H_0 = 70 km s^{-1} Mpc^{-1} cosmology, but this has little effect on our conclusions.

2. THE SIMULATIONS

Our simulations were performed with the parallel TreeSPH code GADGET-2 (Springel 2005), employing the fully conservative formulation (Springel & Hernquist 2002) of smoothed particle hydrodynamics (SPH), which conserves energy and
entropy simultaneously even when smoothing lengths evolve adaptively (see Hernquist 1993; O’Shea et al. 2005). Our simulations account for radiative cooling and incorporate a subresolution model of a multiphase ISM to describe star formation and supernova feedback (Springel & Hernquist 2003). Feedback from supernovae is captured in this subresolution model through an effective equation of state for star-forming gas, enabling us to stably evolve disks with arbitrary gas fractions (see Springel et al. 2005b; Springel & Hernquist 2005; Robertson et al. 2006a, 2006c). This is described by the parameter $q_{\text{gas}}$, which ranges from $q_{\text{gas}} = 0$ for an isothermal gas with effective temperature of $10^4$ K, to $q_{\text{gas}} = 1$ for our full multiphase model with an effective temperature $10^8$ K. We have also compared with a subset of simulations which adopt the star formation feedback prescription from Mihos & Hernquist (1994a, 1994b, 1994c, 1996), in which the ISM is treated as a single-phase isothermal medium and feedback energy is deposited as a kinetic impulse. We examine the effects of these choices in Section 4.3.5, and find they are minimal.

Likewise, although they make little difference to the analysis here, supermassive BHs are usually included at the centers of both progenitor galaxies. These BHs are represented by “sink” particles that accrete gas at a rate $M$ estimated from the local gas density and sound speed using an Eddington-limited prescription based on Bondi–Hoyle–Lyttleton accretion theory. The bolometric luminosity of the BH is taken to be $L_{\text{bol}} = c \epsilon_r M c^2$, where $\epsilon_r = 0.1$ is the radiative efficiency. We assume that a small fraction (typically $\approx 5\%$) of $L_{\text{bol}}$ couples dynamically to the surrounding gas, and that this feedback is injected into the gas as thermal energy, weighted by the SPH smoothing kernel. This fraction is a free parameter, which we determine as in Di Matteo et al. (2005) by matching the observed $M_{\text{BH}}-\sigma$ relation. For now, we do not resolve the small-scale dynamics of the gas in the immediate vicinity of the BH, but assume that the time-averaged accretion rate can be estimated from the gas properties on the scale of our spatial resolution (roughly $\approx 20$ pc, in the best cases). While the BHs can be indirectly important, owing to their feedback ejecting gas into the halo and thus preserving it from star formation until the final merger, we find that, for a given gas content at the time of the actual merger, our results are unchanged in a parallel suite of simulations without BHs.

The progenitor galaxy models are described in Springel et al. (2005b), and we review their properties here. For each simulation, we generate two stable, isolated disk galaxies, each with an extended dark matter halo with a Hernquist (1990) profile, motivated by cosmological simulations (Navarro et al. 1996; Busha et al. 2005), an exponential disk of gas and stars, and (optionally) a bulge. The galaxies have total masses $M_{\text{vir}} = V_{\text{vir}}^3/(10G H(z))$, with the baryonic disk having a mass fraction $m_d = 0.041$, the bulge (when present) having $m_b = 0.0136$, and the rest of the mass in dark matter. The dark matter halos are assigned a concentration parameter scaled as in Robertson et al. (2006c) appropriately for the galaxy mass and redshift following Bullock et al. (2001). We have also varied the concentration in a subset of simulations, and find it has little effect on our conclusions (because the central regions of the galaxy are, in any case, baryon-dominated), insofar as they pertain to disk survival in mergers (it has been demonstrated that halo concentrations are important for, e.g., the exact sizes and velocity scalings of disks, and our predicted disk sizes scale accordingly). The initial disk scale-length is computed based on an assumed spin parameter $\lambda = 0.033$, chosen to be near the mode in the $\lambda$ distribution measured in simulations (Vitvitska et al. 2002), and the scale-length of an initial bulge (when present) is set to 0.2 times this.

Typically, each galaxy initially consists of 168,000 dark matter halo particles, 8000 bulge particles (when present), 40,000 gas and 40,000 stellar disk particles, and one BH particle. We vary the numerical resolution, with many simulations using twice as many particles (and a subset using up to 128 times as many particles). We choose the initial seed mass of the BH either in accord with the observed $M_{\text{BH}}-\sigma$ relation or to be sufficiently small that its presence will not have an immediate dynamical effect, but we have varied the seed mass to identify any systematic dependences. Given the particle numbers employed, the dark matter, gas, and star particles are all of roughly equal mass, and central cusps in the dark matter and bulge are reasonably well resolved.

We consider a series of several hundred simulations of colliding galaxies, described in detail in Robertson et al. (2006b, 2006c) and Cox et al. (2006a, 2006b). We vary the numerical resolution, the orbit of the encounter (disk inclinations, pericenter separation), the masses and structural properties of the merging galaxies, presence or absence of bulges in the progenitor galaxies, initial gas fractions, halo concentrations, the parameters describing star formation and feedback from supernovae and BH growth, and initial BH masses.

The progenitor galaxies have virial velocities $V_{\text{vir}} = 55, 80, 113, 160, 226, 320$, and 500 km s$^{-1}$, and redshifts $z = 0, 2, 3,$ and 6, and the simulations span a range in final spheroid mass $M_{\text{BH}} \sim 10^6-10^{13} M_\odot$, covering essentially the entire range of the observations we consider at all redshifts, and allowing us to identify any systematic dependences in our models. We consider initial disk gas fractions (by mass) of $f_{\text{gas}} = 0.05, 0.1, 0.2, 0.4, 0.6, 0.8,$ and 1.0 for several choices of virial velocities, redshifts, and ISM equations of state. The results described in this paper are based primarily on simulations of equal-mass mergers; however, we examine in Section 4.3.3 how our results scale with mass ratio in mixed encounters, down to mass ratios $\sim 1:10$ or so, below which (as we show) the encounters have little noticeable effect. In detail, the simulations studied there are described in Younger et al. (2008) and constitute a complete subset of permutations of our standard galaxy models with mass ratios uniformly sampling the range 1:1 to 1:8. As in our larger set of 1:1 mergers, at each mass ratio we systematically survey the effects of different absolute galaxy mass, orbital parameters, and disk gas fraction (resulting in a typical $30-40$ simulations spanning the full range of orbital parameters and gas fractions of interest, around each mass ratio $\sim 1:1, 1:2, 1:4, \text{and } 1:8$). We have considered more limited studies of minor mergers where we vary, e.g., the ISM equation of state, redshift, initial disk structural properties; as we find in our studies of these parameters in the larger suite of equal-mass mergers, they make no significant difference to our conclusions here.

Once built, pairs of galaxies are placed on parabolic orbits (motivated by cosmological simulations; see Benson 2005; Khochfar & Burkert 2006) with the spin axis of each disk specified by the angles $\theta$ and $\phi$ in standard spherical coordinates. Table 1 lists the orientations in different representative orbits we have sampled. The particular choice of orbits follows Cox et al. (2006b); there are seven idealized mergers (cases b–h) that represent orientations often seen in the literature (for example, case h, where all the angular momentum vectors of the disks and orbit are initially aligned), the rest (i–p) follow Barnes (1988) by selecting unbiased initial disk orientations according
to the coordinates of two oppositely directed tetrahedrons. These orbits are identical to those considered in various other studies, such as Naab & Burkert (2003). For our series of orbits of strong—pathological orbits (such as the aligned case h above) that vary the pericentric passage distance and the energy of the orbit, described in Robertson et al. (2006a) and Cox et al. (2006b), and find that our estimates are robust to these variations. If we consider the distribution of baryonic mass in $v_{\text{rot}}$, we find a clear segregation between bulge and disk components.

Figure 1 shows this for three simulations with large disks in the remnant. There is clearly a bimodal distribution in $v_{\text{rot}}$, with one component having relatively little rotation (the bulge, with a peak near $v_{\text{rot}} \sim 0$), and one component being largely rotationally supported (the disk, with a peak near $v_{\text{rot}} \sim 1$). There are two stellar populations in these remnants, with a clean division in their rotational support.

We therefore choose to take advantage of our full three-dimensional information in the simulations to easily decompose bulges and disks in a simple, automated fashion. For convenience, let us consider the remnant in cylindrical coordinates $x = (R, \phi)$ where the axis of symmetry ($\hat{z}$) is defined by the net angular momentum vector of the baryonic mass in the relaxed remnant. The effective rotational support of any given stellar or gas particle in the simulation is then

$$\bar{v}_{\text{rot}} = \frac{v_\phi}{v_c(r)},$$

(1)

where $v_c$ is the circular velocity

$$v_c = \sqrt{\frac{GM_{\text{enc}}(r)}{r}},$$

(2)

(here $r$ is the three-dimensional radius from the galaxy center).

Notes. List of disk galaxy orientations for major merger simulations. Columns show: (1) the orbit identification (used to refer to each orbit throughout); (2–3) the initial orientation of disk 1 (in standard spherical coordinates); (4–5) the initial orientation of disk 2; and (6) a brief description of some of the orientations.

![Table 1: Disk Orientations](image-url)

Table 1: Disk Orientations

| Name | $\theta_1$ | $\phi_1$ | $\theta_2$ | $\phi_2$ | Comments |
|------|-----------|---------|-----------|---------|----------|
| b    | 180       | 0       | 0         | 0       | Prograde–retrograde |
| c    | 180       | 0       | 180       | 0       | Both retrograde    |
| d    | 90        | 0       | 0         | 0       | Polar–prograde    |
| e    | 30        | 60      | $-30$     | 45      | “Random” (prograde) |
| f    | 60        | 60      | 150       | 0       | Tilted polar–retrograde |
| g    | 150       | 0       | $-30$     | 45      | Retrograde–“random” |
| h    | 0         | 0       | 71        | 30      | Both prograde      |
| i    | 0         | 0       | 71        | 30      | Barnes orientations |
| j    | $-109$    | 90      | 71        | 90      |                    |
| k    | $-109$    | $-30$   | 71        | $-30$   |                    |
| l    | $-109$    | 30      | 180       | 0       |                    |
| m    | 0         | 0       | 71        | 90      |                    |
| n    | $-109$    | $-30$   | 71        | 30      |                    |
| o    | $-109$    | 30      | 71        | $-30$   |                    |
| p    | $-109$    | 90      | 180       | 0       |                    |

Minor merger orientations

| m000 | 0 | 0 | $-30$ | 45 |
| m030 | 30| 0 | $-30$ | 45 |
| m090 | 90| 0 | $-30$ | 45 |
| m150 | 150|0 | $-30$ | 45 |
| m180 | 180|0 | $-30$ | 45 |

3. THE EXISTENCE OF DISKS IN MAJOR MERGER REMNANTS

Robertson et al. (2006a) and Springel & Hernquist (2005), and subsequently Governato et al. (2007) and Naab et al. (2006) have demonstrated that even major mergers can leave remnants with nonnegligible disk components. Nevertheless, we wish to highlight several properties of these disks first, to establish their existence and nature. Moreover, we wish to ensure that we can robustly identify disks in our merger remnants, before going on to analyze the conditions for their survival. In order to do this, we have considered several methods, which include, e.g., fitting the surface brightness profiles to traditional bulge-disk decompositions, (see Robertson et al. 2006a; Hopkins et al. 2008d), kinematic decompositions based on one and two-dimensional velocity maps (Cox et al. 2006b, L. Hoffman et al. 2009, in preparation), and three-dimensional component fitting. These ultimately give similar results, although, e.g., surface brightness profile fits and velocity decompositions can be considerably dependent on the viewing angle and are not especially robust at separating a small disk (in a bulge-dominated system) from, e.g., other kinematic subcomponents or rotating bulges (a well known observational difficulty; see Balcells et al. 2007; Kormendy et al. 2008; Marinova & Jogee 2007; Barazza et al. 2008, and references therein).

We therefore choose to take advantage of our full three-dimensional information in the simulations to easily decompose bulges and disks in a simple, automated fashion. For convenience, let us consider the remnant in cylindrical coordinates $x = (R, \phi)$ where the axis of symmetry ($\hat{z}$) is defined by the net angular momentum vector of the baryonic mass in the relaxed remnant. The effective rotational support of any given stellar or gas particle in the simulation is then

$$\bar{v}_{\text{rot}} = \frac{v_\phi}{v_c(r)},$$

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where $v_c$ is the circular velocity

$$v_c = \sqrt{\frac{GM_{\text{enc}}(r)}{r}},$$

(2)

(here $r$ is the three-dimensional radius from the galaxy center).

If we consider the distribution of baryonic mass in $v_{\text{rot}}$, we find a clear segregation between bulge and disk components.

Figure 1 shows this for three simulations with large disks in the remnant. There is clearly a bimodal distribution in $v_{\text{rot}}$, with one component having relatively little rotation (the bulge, with a peak near $v_{\text{rot}} \sim 0$), and one component being largely rotationally supported (the disk, with a peak near $v_{\text{rot}} \sim 1$). There are two stellar populations in these remnants, with a clean division in their rotational support.

We can, from this plot alone, estimate a robust disk–bulge mass ratio, from fitting, e.g., the sum of two Gaussian components (disk and bulge) to this distribution (or in a nonparametric sense, by assuming the bulge component has a symmetric rotation distribution about its peak, and mirroring the distribution about that, taking what remains to be disk). Our results are not sensitive to the exact details of our decomposition, but we experiment with a few different methods in order to estimate uncertainties on the bulge–disk decomposition which we refer to below. Again, we have repeated our entire analysis using alternative estimators of the disk-to-bulge mass (direct profile fits and velocity profile decompositions), and find that the same scalings apply in all cases (the uncertainties in the decomposition of a given simulation do, however, increase).

Figure 2 shows one simulation from Figure 1 (S0 major-merger remnant), using this method to decompose the remnant into a stellar bulge and stellar disk. We also show the gas
separately, which can cool and therefore forms an extremely thin disk. The properties are exactly what would be expected for a typical bulge–disk system: the “bulge” is a somewhat flattened ellipse with ellipticity $\epsilon \approx 0.3–0.4$ ($H/R \approx 0.6–0.7$), and is a pressure-supported system, with one-dimensional velocity dispersion $\sigma \sim 120–150$ km s$^{-1}$ (depending on the sightline and slit width) and a rotation velocity $\sim 30–50$ km s$^{-1}$. The resulting rotation parameter of the bulge itself ($(V/\sigma)^2 \sim 0.4$) is typical of reasonably rapidly rotating bulges. It is also compact (as expected), with projected $R_e \sim 1–2$ kpc. The stellar disk is like that of a combined thin–thick disk system, with $H/R \sim 0.15–0.2$, and exhibits a flat rotation curve with $V_{\text{max}} \sim 200$ km s$^{-1}$. The disk is rotationally supported with typical values for a disk of similar mass and overall morphology, $V/\sigma \sim 2–3$, and it is far more extended than the bulge ($R_e \sim 10$ kpc, putting it on the observed disk size–mass relation shown below). The properties of the gas disk are similar, with the obvious exception that, since the gas can cool, it forms a very thin disk ($H/R \lesssim 0.05$).

Figure 3 shows the components of another galaxy (Sbc minor-merger remnant) from Figure 1, in the manner of Figure 2. As expected given the smaller bulge-to-disk ratio in this case, the system is more flattened, with even larger rotational support ($V/\sigma \sim 5–10$ and $H/R \lesssim 0.1$ even in the stellar disk). The bulge is clearly a distinct dispersion-supported component, despite being relatively flattened.

Figure 4 shows the components of the third galaxy (elliptical major-merger remnant) from Figure 1, in the manner of Figure 2. We show this case to demonstrate that embedded disks can be recovered reliably, and are indeed real, rotation supported ($V/\sigma \sim 3–5$), and relatively thin ($H/R \lesssim 0.2$ in the stellar disk, $\lesssim 0.1$ in the gaseous disk) kinematic objects even in simulations where they are not a majority of the mass (here, we find $B/T \sim 0.7$).

As a further check that these are indeed real disks, Figure 5 plots the disks in disk-dominated simulation remnants on the baryonic Tully–Fisher and stellar size–mass relations observed for disks of similar morphology (see Bell & de Jong 2001; McGaugh 2005; Shen et al. 2003; Courteau et al. 2007). For convenience, we will use the estimated bulge-to-disk ratio as a proxy for morphology throughout, with the values as labeled in Figure 5. We take their velocities here from the projected disk...
Figure 2. Bulge (left), stellar disk (middle), and gas (right), in the remnant of a 1:2 mass-ratio major merger with \( \sim 20\% \) gas at the time of the merger, on a typical random orbit (the center panel in Figure 1). From top to bottom, panels show: (a) projected surface density (edge-on view), (b) face-on view, (c) mean edge-on scale height \( H/R \) of the component as a function of circular radius \( R = |x| \) (for this projection), (d) edge-on velocity profile: mean velocity \( v(r) \) (blue) and velocity dispersion \( \sigma(r) \) (red), and (e) rotational support measure \( v/\sigma \).

(A color version of this figure is available in the online journal.)

rotation curves where they are flat, and take \( R_e \) as the projected half-mass radius (note that this is different from the exponential disk scale length \( h \); for a pure exponential disk \( R_e = 1.678 h \), and we convert the observations where necessary accordingly).

For each morphological class, our simulations agree well with the observed Tully–Fisher and size–mass relations. Given our limited sampling of very minor mergers with mass ratios \( \gtrsim 1:8 \), we have only a few simulations with final \( B/T < 0.2 \), but those nevertheless agree (as expected, since they have only been slightly modified from the original disk). We stress that we are not claiming to reproduce the Tully–Fisher or stellar size–mass relation of disks in an a priori manner: our (premerger) disks are constructed, by design, to more or less lie on the observed correlations. What we are saying is that, given progenitor disks that are similar to those observed, disks that form after or survive mergers (even major mergers) will remain on the appropriate correlations for their stellar mass and morphology. In short, when disks do survive mergers, they are “real” disks in the observable sense, not highly flattened bulges or unusual kinematic subcomponents.

4. DISK FORMATION IN MAJOR MERGERS

Clearly, even major mergers can and do produce remnants with significant disks. We therefore ask how these disks form, and whether we can derive some analytic expectation for their masses as a function of progenitor and merger properties.

4.1. Components of the Remnant: Surviving Gas Disks

For simplicity, let us begin with the case of an identical 1:1 mass-ratio merger (we will generalize to arbitrary mass ratios in Section 4.3.3 below). Early in the merger, the galaxies experience a first passage and begin to lose angular momentum to the halos, rapidly coalescing on a timescale of order a couple orbital periods. In the final merger and coalescence of the galaxies, the stars which are initially “cold” (i.e., premerger disks) will scatter and violently relax (Lynden-Bell 1967), forming a de Vaucouleurs-like (1948) quasi-spherical, dispersion-supported profile. Of the gas available at the time of the final merger, some will lose its angular momentum, fall into the galaxy center, and (given the sudden rapid increase in density) rapidly transform into stars in a central starburst, forming a compact, central dissipational component of the remnant bulge (for a detailed study of this component see Hopkins et al. 2008a, 2008d). Gas that is at sufficiently large radii that it cannot efficiently fall in, or gas that for whatever reason cannot efficiently dissipate or lose angular momentum, will rapidly see the central potential relax (the equilibration timescale of the central bulge is only \( \sim 10^8 \) yr) and, having conserved its angular momentum, will rapidly cool.
Figure 3. Bulge, disk, and gas, shown as in Figure 2, for the remnant of a 1:8 mass-ratio minor merger with \( \sim 15\% \) gas, on a polar orbit.

(A color version of this figure is available in the online journal.)

and re-form a thin, rotationally supported disk. Barnes & Hernquist (1996) outline this process and show, in detail, how the cooling gas that survives the merger rapidly settles into a typical, rotationally supported exponential disk. This will then form stars, which constitute a new stellar disk.

We emphasize these three components:

**Premerger stars:** these (along with the dark matter) constitute the collisionless (dissipationless) component of the merger. Because they are collisionless, the stars and dark matter distributions mix in the merger. A given star, as it moves through the merging galaxies on a random orbit, feels a rapidly fluctuating potential, which deflects its orbit and allows for the phase space distribution of the particles to uniformly mix. This violent relaxation process gives rise to a pressure-supported system dominated by random velocities (Lynden-Bell 1967) and transforms initially exponential disk into quasi-spherical de Vaucouleurs-like (1948) S´ersic-law profiles. In the limit of a 1:1 mass-ratio merger, it is a good approximation to assume that all of the stars are violently relaxed—the merger is sufficiently “violent” that no significant component of the premerger stellar disks will survive the merger. This is not necessarily true at lower mass ratios (see Section 4.3.3), but it simplifies our analysis to begin, while we consider such mergers.

The remaining two components of the remnant can be identified with the gas supply available at the time of the final merger.

**Starburst stars:** this is the remnant of a dissipational starburst, triggered in the merger. Some fraction of the gas will efficiently lose its angular momentum in the merger. Because gas can dissipate energy, it will then necessarily rapidly fall into the center of the merging system (essentially free-falling to the center until the collapsing gas becomes self-gravitating; see Hopkins et al. 2008d). Collecting a large gas supply in the center, the result is a rapid, highly concentrated starburst—in gas-rich cases, this is analogous to that observed in, e.g., nearby merging ULIRGs (Soifer et al. 1984a, 1984b; Scoville et al. 1986; Sargent et al. 1987, 1989) and recent merger remnants (Kormendy & Sanders 1992; Hibbard & Yun 1999; Rothberg & Joseph 2004). This builds up a dense, compact central stellar distribution that raises the central phase space density and yields an effectively smaller, more baryon-dominated remnant. The starburst stars, being so concentrated (and typically having a more mixed orbital distribution owing to the random velocities of infalling gas in the starburst), are clearly part of the bulge (although they may have slightly different S´ersic profiles and kinematics from the more extended bulge formed from violent relaxation of the premerger stars). This component is important for the structure and scalings of the bulge/spheroid component, and we study it in detail in Hopkins et al. (2008a, 2008g, 2008d). It is essentially the dissipational component of the merger. For our purposes here, however, this is the gas “lost,” which becomes part of the bulge and no longer contributes to the remnant disk.
Surviving gas/"postmerger" stars: the gas that does not lose its angular momentum will, as described above, form a new disk as the remnant Relaxes. For a 1:1 merger, since (as noted above) the entire stellar distribution is violently relaxed, the postmerger disks can be entirely identified with gas that survives the merger. It is not, in this case, so much that the initial disks survive the merger intact, as it is that some of the gas remains at large radii with significant angular momentum, which can rapidly re-form the disk after the merger. Essentially then (for major mergers), the question of how much of a disk will remain postmerger is a question of how much of the gas (at the time of the merger) will or will not lose its angular momentum.

4.2. How Does the Gas Lose Its Angular Momentum?

How, then, does the gas lose angular momentum in a merger? The basic process has been understood since early simulations involving highly simplified models for gas dissipation in Noguchi (1987, 1988), Hernquist (1989), and Barnes & Hernquist (1991). With improved numerical models, Barnes & Hernquist (1996) followed this process in detail, and showed that what happens in a typical major merger is as follows: the nonaxisymmetric perturbation (owing to the companion) in the system induces (largely after first passage and on the final coalescence, since this is where the interaction is significant) a nonaxisymmetric response in the disk. A stellar bar and gas bar form, but because the gas is collisional and the stars are collisionless, the stellar bar will trail or lag behind the gas bar by a small offset (typically ~a few degrees). The stellar bar therefore torques the gas bar, draining its angular momentum, and causing the gas to collapse to the center.

Figure 6 illustrates this in a couple of representative 1:3 mass-ratio major merger simulations. For a more detailed description and illustration of the relevant physics, we refer to Barnes & Hernquist (1996, particularly their Figures 3–8); but we briefly outline the scenario here. We show the morphology of the gas before the merger, when the disk is undisturbed, and shortly after both first and second passages (the second passage leading, in these cases, to a rapid coalescence), as well as in the relaxed remnant. The barlike nonaxisymmetric perturbation induced by the close passages is clearly evident; the stars show a similar morphology at each time, with a small phase offset in the bar. In what follows, we will refer to this nonaxisymmetric response as a "bar," for simplicity and because morphologically the induced feature, at least for some time during the merger resembles bars in isolated barred spirals. However, we caution that the formation mechanism which excites this response may be different from that causing bars in isolated galaxies. Furthermore, while the nonaxisymmetry is present throughout the merger, it at times would not be classified as a bar morphologically, particularly during the final coalescence of the galaxy nuclei, when the resulting gas inflows are strong.
pattern and (in the remnant) a stellar bulge. As we discuss in Section 4.3.2, a prograde encounter, being in resonance, produces a noticeably more pronounced bar distortion (both in amplitude—effectively “bar mass”—and spatial extent). Shortly after each passage, this double bar system efficiently removes angular momentum from the gas, allowing it to fall into the center of the galaxy and participate in a centrally concentrated starburst.

Following Barnes & Hernquist (1996), we track the gas in the primary disk that will turn into stars in the final starburst, calculating the net instantaneous gravitational torque decelerating the disk rotation. We can coarsely infer what the total effective torque must be by simply differentiating the specific angular momentum of this gas at a given time, and compare this to the net torque from different sources. Specifically, we separate the instantaneous gravitational torques into the internal torques—those from the stellar disk in the same galaxy as the gas, chiefly from the bar (since the axisymmetric disk, by definition, exerts no net torque), and the external torques—those from the gravity of the secondary galaxy itself and the extended halos and their substructure.

It is clear that, especially for the phases of interest shortly after second passage and leading into the final starburst, when this gas loses its angular momentum, the total torques are dominated by internal torques from the stellar disk/bar system. The agreement between these torques and the rate of change in the specific angular momentum further argues that there are no other major sources of angular momentum loss (specifically, both this comparison and direct calculation demonstrate that the “hydrodynamic torques” defined by pressure forces are not dominant).

As a result of these torques, gas within some critical radius where the internal torques are strong (roughly inside the “bar radius” in Figure 6) rapidly loses angular momentum. We define this radius more precisely in Section 4.3, but it is clear in the figure that at sufficiently large radius the bar perturbation is weaker (and moreover, at larger radius the potential of the disk, whether barred or unbarred, appears increasingly axisymmetric); gas outside of these radii is relatively unaffected. In general, then, the means for a more efficient encounter to consume a larger fraction of the gas in the disk is to induce a stronger bar disturbance, which is able to effectively exert internal torques out to larger radii, stripping more gas of angular momentum and bringing it into the central starburst (as evident in the stronger prograde encounter in Figure 6). Finally, the system relaxes—the gas that has not been subjected to strong internal torques, having retained its angular momentum (at least in large part), can rapidly re-form a disk. This may entail some redistribution of that angular momentum (“filling in” where the bar depleted the gas of the disk), but does not lead to further significant angular momentum loss.

Barnes & Hernquist (1996) and Barnes (1998) illustrate that this internal torquing is by far the dominant source of angular momentum loss, for typical orbits. This is because the stellar bar is: (1) more or less aligned in the plane with the gas bar, (2) trailing it by a small amount, and (3) relatively long-lived (it lives the rest of the duration of the merger, as opposed to the short time that is, e.g., pericentric passage). The companion itself (either its...
baryonic mass or its halo), in most orbits, is not perfectly aligned with the gas disk, and the torque directly from it is much weaker (the tidal torquing drops by a factor $\sim (R_{disk}/R_{peri})^3$), and it can act only for a short duration on pericentric passage. There are some pathological orbits (e.g., perfectly coplanar prograde orbits) where this is not true, but these are exceptional cases, and we discuss them in Section 4.3.6.

At the final merger, one might imagine that mixing of random gas orbits or collisions and shocks would rapidly drain angular momentum, similar to what happens to the stars in violent relaxation. However, this is not possible, precisely because the gas is collisional: a Lagrangian gas element cannot go back and forth through the galaxy, but sticks to the other gas which has some net angular momentum. There could in principle be some net angular momentum cancellation, but this is inefficient—the net angular momentum will almost always be comparable to the initial total. Even assuming random cancellation between two disks with comparable absolute angular momentum, the average change in net specific angular momentum is a factor of $\sim 2/3$; when one accounts for the angular momentum of the orbit—typically comparable or even larger than that in the disks—there is often no change or even a net gain in the gas specific angular momentum in a merger.

A proper calculation shows that over the range in mass ratios $\mu \sim 0.1$–1, for a range of typical impact parameters $b \sim 0.5$–$5 R_d$, the expected final specific angular momentum after cancellation is approximately equal to the initial specific angular momentum of the primary (with $\sim 20\%$ scatter). Cancellation is therefore inefficient. Even these cancellations, we find in detail, do not generally lead to a starburst in the same manner as a merger-induced bar, but simply lead to moderate disk contraction (and an equal number of mergers will scatter toward the opposite sense leading to disk expansion, keeping a mean specific angular momentum that is constant). They do

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**Figure 6.** Illustration of the key processes that drive starbursts in a merger. Top: projected gas density (as in Figure 2) in the plane of the disk at representative times in a retrograde l:3 merger (left to right: before interaction, just after first passage, just after second passage/coalescence, after relaxation). For clarity, just the gas from the primary is shown. Middle: same, for a prograde encounter. The passage of the secondary induces a barlike nonaxisymmetric disturbance in the primary, which survives after the short-lived passage and removes angular momentum from the gas, leading to a starburst. The same process occurs on both passages, with a larger (albeit less “barlike”) asymmetry on coalescence. The prograde encounters, being in resonance, induce a stronger response that extends to larger radii. Bottom: quantities of interest in the merger. Left: star formation rate (SFR) as a function of time in the prograde encounter (solid line; retrograde is similar, but with a weaker enhancement in bursts). The red dotted line shows the specific angular momentum of the gas that will participate in either burst, in arbitrary units; as the gas rapidly loses angular momentum after the passages, it drives a central starburst. Center left: net specific torque on the primary gas that will participate in the final, central starburst, as a function of time in the merger (in units of the initial total angular momentum per Gyr). We compare the roughly numerically estimated net torque (diamonds; from differentiating the specific angular momentum of the gas) and the torque from two sources: stars in the same disk as the gas (internal torques; black thick line), and the secondary galaxy and extended halos (external torques; red dot-dashed line). The loss of angular momentum that drives the secondary burst at $t \sim 2.2$–2.5 is driven by internal torques from the disturbed stellar disk; not the torque from the secondary galaxy itself. Center right: final specific angular momentum content of material (gas plus stars) that was originally gas at $R < R_{crit}$ (our predicted radius where merger-induced internal torques should be efficient at removing angular momentum) or at $R > R_{crit}$. There is a strong division: gas inside a characteristic radius (corresponding to where the internal asymmetry is strong; akin to the corotation resonance) is mostly stripped of angular momentum. Gas at larger radii conserves sufficient angular momentum to maintain a disk at similar $R_d$ and $V_d$. Right: original (cumulative) radial distribution of gas that participates in the final starburst, relative to the initial disk effective radius (same in all cases) for the prograde and retrograde cases shown and a more minor retrograde merger. More resonant (prograde) and more major encounters induce a stronger response, with a larger corotation radius, and so torque gas out to larger radii and efficiently strip more gas of angular momentum as predicted.

(A color version of this figure is available in the online journal.)
not cause a starburst because, if two random parcels or streams of gas shock and lose angular momentum, the alignment and relative momenta would have to be near-perfect for them to lose, say 95% of their angular momentum and fall all the way to the central ∼ 100 pc where a nuclear starburst would occur. Rather, they will lose some fraction of order unity of their angular momentum, fall in to a slightly smaller radius, and continue to orbit.

Without the bar that can continuously drain angular momentum, the true burst is indeed inefficient. This initial bar-induced angular momentum loss scenario has been well established in subsequent numerical studies (see Noguchi 1987, 1988; Hernquist 1989; Hernquist & Barnes 1991; Borderies et al. 1989; Barnes & Hernquist 1991, 1996; Barnes 1998; Mihos & Hernquist 1996; Springel & Hernquist 2005; Robertson et al. 2006a; Cox et al. 2006b, 2008; Berentzen et al. 2007; Naab et al. 2006). We therefore can simplify our question to ask: how efficient will a given lagging stellar bar be at removing angular momentum from a leading gas bar?

4.3. A Simple Model: Overview

Consider a disk that contains a total gravitational mass $M$ (which can include a bulge and dark matter as well; the disk mass fraction we will denote $f_{\text{disk}}$) within a characteristic scale length $R_d$. Some convenient dimensional variables are

\[
M_{\text{disk}} = f_{\text{disk}} M \quad (\text{disk mass})
\]

\[
M_{\text{bar}} = f_{\text{bar}} M \quad (\text{stellar bar mass})
\]

\[
v_c = \sqrt{GM/R_d} \quad (\text{characteristic circular velocity})
\]

\[
\Omega_d = v_c/R_d \quad (\text{characteristic frequency})
\]

\[
P = \frac{2\pi}{\Omega_d} \quad (\text{rotation period}).
\]

We also define the disk thickness according to a characteristic (assumed constant) relative scale height (height $H$ versus radius $R$; $H/R =$ constant)

\[
\tilde{H} \equiv H/R \quad (\text{disk scale height}).
\]

In these units, the circular velocity at a given cylindrical radius $R$ is given by

\[
v_{\text{circ}}(r) = v_c \tilde{v}(r) \equiv v_c \sqrt{M_{\text{enc}}(r)/M} \frac{R_d}{R}
\]

and dimensionless lengths are defined by

\[
\tilde{x} = x/R_d
\]

\[
\tilde{y} = y/R_d
\]

\[
\tilde{R} \equiv \sqrt{x^2 + y^2}/R_d.
\]

Throughout, we will use this notation: e.g., the dimensional variable $u$ is equal to the dimensionless variable $\tilde{u}$ times the appropriate combination of dimensional constants above.

We will show that, in such a disk, a gas bar with a lagging stellar bar will efficiently cause gas to lose its angular momentum and dissipate into a central starburst. This will be the case for gas interior to a radius

\[
\frac{R_{\text{gas}}}{R_d} \leq 1 - \alpha (1 - f_{\text{gas}}) f_{\text{disk}} F(\theta, b) G(\mu),
\]

where $\alpha \sim 1$ is an appropriate integral constant (depending weakly on details of the stellar profile shape and bar dynamics), $f_{\text{gas}}$ is the gas fraction in the disk and $f_{\text{disk}}$ is the disk mass fraction. The factor

\[
G(\mu) = \frac{2\mu}{(1 + \mu)}
\]

contains the dependence on the merger mass ratio (where $\mu \leq 1 \equiv M_2/M_1$). The term

\[
F(\theta, b) \equiv \left( \frac{1}{1 + (b/R_d)^2} \right)^{\frac{3}{2}} \frac{1}{1 - \Omega_c/\Omega_d}
\]

accounts for the orbital parameters: $b$ is the distance of pericentric passage on the relevant final passage before coalescence ($\sim 1$--a couple $R_d$, for typical cosmological mergers) and $\Omega_c$ is the orbital frequency at pericentric passage,

\[
\frac{\Omega_o}{\Omega_d} = \frac{v_{\text{peri}} R_d}{v_c} \frac{R_d}{b} \cos(\theta)
\]

\[
= \sqrt{2(1 + \mu)[1 + (b/R_d)^2]^{-\frac{3}{4}}} \cos(\theta)
\]

\[
\approx 0.6 \cos(\theta),
\]

where $\theta$ is the inclination of the orbit relative to the disk, and the last equality comes from adopting typical cosmological orbits and mass ratios (but in any case, this is quite weakly dependent on the mass ratio).

In the following sections, we derive this scaling piece by piece, and compare each aspect to the results from our library of hydrodynamic simulations. We show that it is robust and accurate as an approximation to the behavior in full numerical hydrodynamic experiments over a wide dynamic range of several orders of magnitude in surviving disk fraction (from systems with ~ 80%--100% of their disks surviving a merger to systems with < 1% disk after a merger), as well as the entire dynamic range in mass, gas content, orbital properties, and different feedback prescriptions with which we experiment.

4.3.1. Dependence on Disk Gas Content

Let us consider an infinitely thin gas bar (a good approximation, owing to the efficiency of gas cooling) in a potential that is otherwise cylindrically symmetric except for the presence of a stellar bar. For simplicity, assume that the gas bar follows a fixed pattern speed $\Omega_o$ ($\sim \Omega_d$; we will derive the pattern speed later) in the disk (while there is not exactly a constant pattern speed in the outer regions of the disk, the torques are weak there in either case, and this approximation is globally quite good). We take $z = 0$ to be the plane of the disk, and, without loss of generality, consider a frame rotating with the pattern speed of the gas bar, so that the bar lies along the $x$ axis. The material in the bar is rotationally supported, so it has instantaneous velocity $\vec{v}_{\phi} = -v_c \hat{\nu}(R) \hat{y}$, where $v_c \hat{\nu}(R)$ (defined above) is the circular velocity at each point $x$.

Now, consider a stellar bar of total mass $M_{\text{bar}}$ also at fixed pattern speed, but offset by some instantaneous angle $\phi_b$ from the gas bar (i.e., along the axis $y = \tan(\phi_b) x$). The mass per unit length in the bar at a distance $R$ along the bar is

\[
dM_{\text{bar}}/dR = (M_{\text{bar}}/R_d) \tilde{\Sigma} R/R_d,
\]

where $\tilde{\Sigma}$ is the appropriate dimensionless mass profile and $R_d$ is some characteristic scale length (usually corresponding closely to the scale length of the unperturbed disk). If the initial disk is in equilibrium (i.e., if the bar is some reasonable perturbation to the initial system),
then the unperturbed net acceleration in the $x$ direction at some point $x$ in the gas bar will just be canceled by the rotation of the system. Of interest here is the torque; if the stellar bar is also thin, then at a point $x = \tilde{x} R_d$ in the gas bar, the net torque per unit mass from the stellar bar will be

$$\frac{dj}{dt} = \tilde{x} R_d \frac{dv_y}{dt} = -G \frac{M_{\text{bar}}}{R_d} I_0(\phi_b, \tilde{x}),$$

(11)

where $I_0 \sim 1$ is a dimensionless integral which depends weakly on $\phi_b$ and $\tilde{x}$ (at large $\phi_b$, $I_0 \to 0$, reflecting the fact that the torque is dominated by times when the bars are close; since $\phi_b \ll 1$ is expected, it is a good approximation to ignore the $\phi_b$ dependence of $I_0$).

If we assume that the stellar bar is infinitely thin, there is a weak divergence in $I_0$ as $\phi_b \to 0$ (the accelerations become large when the bars nearly overlap). More accurately, we can allow for some finite height in the stellar disk/bar (it will always be thicker than the gas disk/bar). Let the stellar bar have a constant relative scale height $H/R$ given by Equation (4), and for simplicity take its vertical profile to be constant density out to a height $\pm H$ (although assuming a more realistic vertical profile $\exp(-|z|/H)$ or $\sech^2(z/H)$ makes almost no difference to our calculation). The specific torque at $x$ in the gas bar now becomes

$$\frac{dj}{dt} = -G \frac{M_{\text{bar}}}{R_d} \frac{1}{\sqrt{\sin^2 \phi_b + \frac{H^2}{R_d^2}}} I_1(\phi_b, \tilde{x}, \tilde{H}),$$

(12)

where $I_1$ is an even weaker function of $\phi_b$ and $\tilde{x}$ than $I_0$. The important behavior is entirely captured ignoring $I_1$, namely that the finite width of the stellar bar suppresses the numerical divergence seen earlier.

The bar mass $M_{\text{bar}}$ represents the stellar mass in the disk that is effectively part of the bar at the appropriate instant. We can therefore parameterize $f_{\text{bar}} = M_{\text{bar}}/M$ as

$$f_{\text{bar}} = (1 - f_{\text{gas}}) f_{\text{disk}} \Psi_{\text{bar}}.$$  

(13)

Here, $f_{\text{disk}}$ is the disk mass fraction, and $f_{\text{gas}}$ is the gas fraction in the disk (since we are interested in the cold gas, we explicitly ignore gas in, e.g., a bulge or hot halo component). Therefore, the stellar mass of the disk is $(1 - f_{\text{gas}}) f_{\text{disk}}$—this defines the maximum mass that could be in the stellar bar. The parameter $\Psi_{\text{bar}}$ thus defines the bar “efficiency”—in an instantaneous sense as we have defined it here, $\Psi_{\text{bar}} = 0$ means there is no stellar bar, $\Psi_{\text{bar}} = 1$ implies the maximal stellar bar.

Already, we have one significant scaling—the bar strength, and correspondingly the strength of the torques on the gas, scale with $(1 - f_{\text{gas}})$. In very gas-rich systems where $f_{\text{gas}} \to 1$, there is no stellar mass to form a lagging bar and remove angular momentum from the gas. The gas itself may form a bar, but without a stellar bar to drag it, the angular momentum loss (over the timescales of relevance for a merger) is inefficient. This is well known in, e.g., dynamical studies of pure gas and stellar bars (e.g., Schwarzschild 1981; Athanassoula et al. 1983; Pfenniger 1984; Combes et al. 1990; Friedli & Benz 1993; O’Neill & Dubinski 2003). There might be some angular momentum loss in such a case between, e.g., bar and halo (e.g., Hernquist & Weinberg 1992a), but it will be small—certainly nowhere near

the efficient stripping of angular momentum needed to induce a significant starburst.

If we consider a Lagrangian gas element at some initial radius $x_0$, then its orbit will decay as it loses angular momentum. The instantaneous rate of change in the radius $R = |x|$ will be given by $dR/dt = -(R/v_\phi) dv_y/dt$. The characteristic timescale for the system to evolve is given by $2\pi/\Omega_d$, where $\Omega_d \sim v_\phi/h$ is the characteristic frequency of the disk. Define the timescale

$$\tau \equiv \frac{\Omega_d}{2\pi} \frac{t}{v_c} = \frac{\sqrt{G M}}{R_d} \sqrt{\frac{H}{v_c}},

(14)

where $M$ is the total effective gravitational mass of the disk and $R_d$ is again a characteristic scale length. We now have

$$\frac{d\tilde{x}}{d\tau} = \frac{2\pi}{\Omega_d} \frac{1}{v_c} \frac{G M_{\text{bar}}}{R_d^2} I_1 \equiv 2\pi f_{\text{bar}} I_2(\phi_b, \tilde{x}).$$

(15)

Because the merger occurs on a couple of dynamical timescales, i.e., a time $\Delta \tau \sim 1$, to lowest order (ignoring, e.g., the complications of different orbital parameters) we expect that gas within a radius

$$\tilde{x} \ll \Delta \tau \frac{d\tilde{x}}{d\tau}$$

(16)

will efficiently lose angular momentum and fall to the center, becoming part of the central starburst, while gas at $\tilde{x} \gg \Delta \tau d\tilde{x}/d\tau$ will avoid the starburst. This defines a scale

$$\frac{R_{\text{gas}}}{R_d} \lesssim (1 - f_{\text{gas}}) f_{\text{disk}} \Psi_{\text{bar}}(\phi_b, \tilde{H}, \ldots)$$

(17)

within which the gas will lose angular momentum. For convenience, we have collected all of the dimensionless integral factors, including $\Psi_{\text{bar}}$ (the efficiency of forming the stellar bar) and the dynamical integral factors (e.g., $I_1, I_2$) from above, into the term $\Psi_{\text{bar}}$ that represents the full solution. We write $\Psi_{\text{bar}}(\phi_b, \tilde{H}, \ldots)$ because, as we will show, this quantity (at present) encapsulates our ignorance of, e.g., the orbital parameters and merger mass ratio; for a 1:1 merger on a typical orbit, however, $\Psi_{\text{bar}} \sim 1$.

The total gas mass which will lose angular momentum and fall into the center of the galaxy will be $f_{\text{gas}} \times f(< R_{\text{gas}})$, where $f(< R_{\text{gas}})$ is the mass fraction within the characteristic radius above, according to the details of the mass profiles and dimensionless integral above. We consider solutions for a variety of profiles, including, e.g., an exponential, isothermal sphere, and a Mestel (1963) $1/R$ disk profile. In general, we find that there is little difference between the predictions for these various profiles—the differences in the mass profile shapes tend to cancel out and leave only weak corrections to the simple dimensional scaling. An exponential disk with $\Sigma \propto \exp(-R/R_d)$ contains a mass fraction

$$1 - (1 + R/R_d) \exp(-R/R_d),$$

(18)

within a radius $R$; here $\Psi_{\text{bar}}$ must be solved numerically. We obtain nearly identical predictions, however, assuming a 1/$R$ disk or an isothermal sphere profile for the gas, which allows us to analytically solve the relevant equations and write the predicted gas fraction consumed in the form

$$f_{\text{burst}} = f_{\text{gas}} (1 - f_{\text{gas}}) f_{\text{disk}} \Psi_{\text{bar}}(\phi_b, \tilde{H}, \ldots).$$

(19)
where \( \Psi_{\text{bar}} \approx 1 \) can be analytically calculated for these profiles (with the equations above) under certain conditions: if the dependence on the orbital parameters is separable and we define \( \Psi_{\text{bar}} \) by the requirement that the radius in Equation (17) satisfy \( \tilde{x} = \Delta \tau \frac{d \tilde{l}}{d \tau} \), then for instantaneous bar lag of \( \phi_b \) a few degrees in a thick disk of height \( \tilde{H} \sim 0.2 \), \( \Psi_{\text{bar}} \) is given by

\[
\Psi_{\text{bar}} \approx F(\ldots) \left[ 1 - \exp \left( -\frac{\sin(2 \phi_b)}{\sin^2(\phi_b) + H^2} \right) \right] \sim 1,
\]

where we explicitly show \( F(\ldots) \) as we have suppressed our ignorance of the orbital parameters. Nevertheless, this simple scaling alone provides a remarkably successful description of many of our simulations.

Figure 7 tests this simple prediction. For a suite of merger simulations, we compare the mass fraction in the central starburst, \( f_{\text{burst}} \), to the gas content of the (immediately premerger) disks, \( f_{\text{gas}} \). We can either determine the starburst mass fraction by directly measuring the gas mass that loses its angular momentum and participates in the brief nuclear starburst, or by measuring the gas content that survives and forms a disk (described in Section 4) and assuming the gas that did not survive (relative to that available just before the final merger) was part of the burst. In either case, we obtain a nearly identical answer for each simulation. For now, we consider only simulations with a 1:1 mass ratio—we will generalize to arbitrary mass ratios below. In all these simulations, the premerger \( f_{\text{disk}} \approx 1 \), and measuring \( \phi_b \) and \( \tilde{H} \) just before the merger we expect \( \Psi_{\text{bar}} \approx 1 \). We consider one set of orbits at a time—i.e., compare only systems with the same orbital parameters, so that we can temporarily suppress the dependence on them (this yields a systematic offset between each set of orbital parameters—the solutions plotted account for that following our solution in the next section). At a fixed orbit, for these mergers, then, the only parameter that matters should be \( f_{\text{gas}} \). The simulations at each orbit span a wide range in \( f_{\text{gas}} \), from \( \sim 0.01 \) to 1.

We compare the relation between \( f_{\text{burst}} \) and \( f_{\text{gas}} \) resulting from the full numerical experiments to the simple scalings predicted by Equations (17)–(19). In detail, we show two solutions—first, the scaling given by Equation (19), \( f_{\text{burst}} \propto f_{\text{gas}}(1 - f_{\text{gas}}) \), appropriate for an isothermal sphere or Mestel (1963) disk profile with \( \Phi_{\text{bar}} \approx 1 \); and second, the appropriate numerical solution (following Equations (17)–(18)) for an exponential disk. In either case the analytic solutions are similar, and agree well with the trend seen in the simulations. It is clear that the efficiency of the burst in simulations is—as we predict—not constant. It is not the case that the entire gas supply is always stripped of angular momentum and consumed in the final merger (which would yield \( f_{\text{burst}} = f_{\text{gas}} \)). Rather, when \( f_{\text{gas}} \) is sufficiently high, only a fraction \( \sim (1 - f_{\text{gas}}) \) of the available gas is able to efficiently lose its angular momentum and participate in the starburst.

Figure 8 repeats this comparison in terms of the surviving disk mass. We argued that the gas that does not lose angular momentum in the merger will survive to re-form a disk. Because these are 1:1 mergers where we can safely assume that the entire stellar disks are destroyed, we expect then that the disk mass fraction will be \( f_{\text{disk}} = f_{\text{gas}} - f_{\text{burst}} \). Using the method described in Section 4 to estimate the remnant disk mass fractions, we plot \( f_{\text{disk}} \) versus \( f_{\text{gas}} \) for each of several orbital parameter sets. Again, the exact details of the predictions depend on orbital parameters in a manner we derive below, but for now we are interested in whether or not they obey the predicted scaling with \( f_{\text{gas}} \). Indeed, they do. Over 2 to 3 orders of magnitude in fractional disk mass (and \( \sim 5–6 \) in absolute disk mass), the simple scaling here agrees well with full numerical experiments. It is clear that some of \( f_{\text{gas}} \) is consumed, as expected (if all the gas survived, we would obtain \( f_{\text{disk}} = f_{\text{gas}} \)), but in fact, especially at low \( f_{\text{gas}} \), the efficiency of angular momentum loss is high as predicted and the gas participates in the starburst.

Figure 9 simplifies this—we again compare \( f_{\text{disk}} \) and \( f_{\text{gas}} \), but effectively put all the orbits on the same footing by plotting \( f_{\text{disk}} \) versus the predicted \( f_{\text{disk}}(f_{\text{gas}}, \ldots) \) (i.e., including the orbital parameters according to the predictions in Section 4.3.2 below). Essentially this amounts to implicitly including \( F(\ldots) \) in Equation (19) above. The remaining scaling should just represent the predicted \( f_{\text{burst}} \propto f_{\text{gas}}(1 - f_{\text{gas}}) \). For each simulation, we show an error bar corresponding to the range of \( f_{\text{disk}} \) estimated using different methods (e.g., a full three-dimensional kinematic decomposition, one and two-dimensional kinematic modeling, and surface brightness profile fits, as described in Section 4). The agreement is surprisingly good, given the simplicity of our derivation. Moreover, the scatter is quite small—a factor \( \sim 2 \) to 3

\[\text{Figure 7. Mass fraction formed in the central, dissipational starburst as a function of gas mass fraction at the time just before the starburst, in a suite of major 1:1 mass-ratio mergers. The solid lines are our theoretical predictions (in the equations above). The dotted line corresponds to busting all the available gas (in the equations above). We show results here for several orbits from Table 1: a random orbit with both disks inclined (e: top), an inclined polar–prograde orbit (k: middle), and a polar–polar orbit (f: bottom). The simulations agree well with our analytic predictions: more gas-rich mergers are less efficient at torquing angular momentum away from the gas and funneling it into the starburst (efficiency } \sim (1 - f_{\text{gas}})).\]

\[\text{(A color version of this figure is available in the online journal.)}\]
Figure 8. Relaxed postmerger remnant disk mass fraction vs. gas fraction just before the merger, for 1:1 major mass-ratio mergers. In this case essentially all the premerger stellar mass is transformed (violently relaxed) into bulge—the disk is formed from the gas that survives the merger. Panels consider different orbits, with points as Figure 7. The solid lines are our theoretical predictions \( f_{\text{disk}} = f_{\text{gas}} \{1 - (1 - f_{\text{gas}}) \Psi \} \), see Equation (19), the dotted lines correspond to all the gas surviving and forming a disk \( f_{\text{disk}} = f_{\text{gas}} \). Again, the simulations agree well with our analytic predictions; gas-rich mergers are inefficient at stripping angular momentum from the gas, leaving significant gas content that rapidly re-forms a postmerger disk.

(A color version of this figure is available in the online journal.)

Figure 9. Relaxed postmerger remnant disk mass fraction vs. our analytic predictions as a function of gas fraction and orbital parameters, for 1:1 mass-ratio mergers. Error bars correspond to variation using different methods to estimate the disk–bulge decomposition.

(A color version of this figure is available in the online journal.)

at very low \( f_{\text{gas}} \) and considerably smaller \( \lesssim 50\% \) at high \( f_{\text{gas}} \). It seems that our simple scaling indeed captures the most important physics of angular momentum loss—namely that with less fractional stellar material in the disk, there is less mass available to torque on the gas bar in a merger, therefore less angular momentum loss in the gas.

4.3.2. Dependence on Orbital Parameters

We now turn to how the details of the orbit affect the loss of angular momentum. Before, we made the simplifying assumptions that the stellar bar lagged by some constant angle \( \phi_b \) and that the characteristic time for the perturbation to act was of the order of the disk rotational period/dynamical time. While these turn out to give reasonable scalings, we can improve upon them.

The secondary (“perturber”) galaxy will have some characteristic orbital frequency \( \Omega_o \), approximately given by

\[
\Omega_e \sim \frac{v_{\text{per}}}{b},
\]

where \( v_{\text{per}} \) is the velocity at pericentric passage and \( b \) is the impact parameter or pericentric passage distance. Because the behavior we are interested in is relatively short-lived, this is reasonable even for first passages or “flyby” encounters—we are interested in the orbital frequency at pericentric passage because this is when the forcing is strongest and the bar is driven.

Now, consider the frame rotating with the disk/bar at a frequency \( \sim \Omega_d \). In this frame, the secondary galaxy will have an apparent frequency for an orbit projected into the disk plane of

\[
\Omega_{\text{eff}} = \Omega_o \cos \theta - \Omega_d,
\]

where \( \theta \) is the inclination of the orbit relative to the plane of the disk (in standard parlance for orbital parameters as described in Section 2). Note that we are interested in the component of the orbital motion in the plane of the disk: in terms of the standard orbital parameters \( \phi \) and \( \theta \) of the primary galaxy (angle of the angular momentum vector of the disk relative to the plane of the orbit), this is \( \Omega_o \cos \theta \). For a case with, e.g., \( \theta = 0 \) (prograde) and a parabolic orbit with small impact parameter (\( \Omega_e \sim \Omega_o \)), the system is maximally prograde—in the frame of the rotating disk there is almost no net circular motion of the secondary. For the same orbit but \( \theta = 180^\circ \) (retrograde), the secondary completes a circular orbit around the disk in just half a disk dynamical time (\( \Omega_{\text{eff}} \sim -2 \Omega_o \)). The time required for the secondary to complete a revolution in this frame is therefore (in our dimensionless units \( \tau = t/(2\pi/\Omega_d) \))

\[
\tau_{\text{circ}} = \frac{1}{1 - \frac{\Omega_o}{\Omega_d} \cos \theta}.
\]
The timescale for a gas galaxy to lose its angular momentum and fall to the center of the galaxy is given by our earlier estimate of the torque, as $\tau_{\text{loss}} \sim \tilde{x} / (d\tilde{x}/d\tau)$. If $\tau_{\text{loss}} \ll \tau_{\text{circ}}$ at a given radius, then the derivation we have obtained is essentially valid: the system sees a quasi-static perturbation to the potential, loses its angular momentum, and collapses before the perturbation can damp out or circularize. However, if $\tau_{\text{loss}} \gg \tau_{\text{circ}}$, then the system has not lost much angular momentum by the time the secondary completes a revolution, and will gain some of those losses back as the system comes around the other side. In the limit where $\tau_{\text{circ}}$ is short (much shorter than the local dynamical time), for example, then the potential is effectively circularized—the gas at these radii may undergo oscillatory motion and even have, e.g., spiral waves driven by this external forcing, but there is no means by which the system can introduce a strong net asymmetry to drive inflows.

We can therefore improve our previous estimate: instead of taking $\Delta T \sim 1$ (i.e., a disk rotation period) as the only characteristic timescale, we argue that gas with

$$\frac{\tilde{x}}{d\tilde{x}/d\tau} \lesssim \tau_{\text{circ}} = \frac{1}{1 - \frac{\Omega_{\text{b}}}{\Omega_{\text{d}}} \cos \theta}$$

(24)

will lose its angular momentum, while gas at larger radii will not.

In detail, we can integrate the equations from 4.3.1 for a parcel of gas at some initial radius $x_0$, in a time-dependent potential of this nature. For simplicity, we assume that the secondary drives a circular perturbation in the potential with frequency $\omega = \Omega_{\text{b}}$ and calculate the bar response using the gaseous disk (assuming, again, that it is infinitely cold) and stellar disk (assuming that the scale height $H$ translates into a corresponding velocity dispersion $\sigma / v_c$) wave dispersion relations from Binney & Tremaine (1987). In practice, we find that this is not much different from assuming that the lag in the stellar bar grows with time $\propto \tau_{\text{circ}}$ (i.e., that the stellar bar can keep up or reverse sense tracking the perturbation without significant energy loss, or, more or less equivalently, that the two bars are only in phase when the perturbation is strong, and then rapidly fall out of phase—at the $\gtrsim 5^\circ$–$10^\circ$ level, once the perturbation is weak or reverses its sense as the phase of the secondary reaches $\gtrsim \pi/2$).

In principle, we now have a physically motivated and fully time-dependent model for $\phi_\rho(t)$ and the response of the gas bar. This allows us to properly integrate out the dependence of $\Psi_{\text{bar}}$ on $\phi_b$ and instantaneous conditions and replace it with the appropriate integral dependence on orbital parameters and disk gas content and structure.

We find that there is a strong division in expected behavior, at more or less exactly the characteristic radius implied by Equation (24). Within this radius, gas (in our simple numerical calculations) is effectively torqued efficiently as it enters the gas bar near resonance (but slightly leading the stellar bar), and plunges to the center. Gas outside this radius begins to feel a perturbation, but then the phase of the secondary cycles around and the sense of the torques begin to weaken or reverse (depending on the details of the orbit), and generate wave motion in the gas but no significant angular momentum loss or infall. Not only is the transition between these two regimes predicted by the simple scalings above, but we find in more detailed numerical calculations that the width of the transition region (where behavior is more sensitive to the details of, e.g., the profile shapes and assumptions about the bars) is quite narrow, $\sim 20\%$ of that radius.

This should not be surprising. Essentially, what we have derived is a rough equivalent of the corotation condition, but for forced bars as opposed to isolated self-generating (swing-amplified) bar instabilities. It is well known from studies of idealized bars (see Schwarz 1981; Pfenniger 1984; Noguchi 1988; Binney & Tremaine 1987; Berentzen et al. 2007) that gas can be efficiently torqued inward of the corotation resonance (in the language above, given the forcing with pattern speed $\Omega_\text{s}$, this is interior to the radius where the relative motion of the secondary is slow relative to the dynamical time, and so the perturbation does not circularize). Moreover, the resonant structure around these radii is known to be sharp; if we follow a derivation similar to Borderies et al. (1989), their derivation is intended to apply to planetary disks with satellites, but the relevant physics is similar) it is straightforward to show that the detailed numerical prefactors will be swamped for all but a narrow range of radii around this resonance by the strong dependence of the resonant forcing on radius (roughly going as some large power of $(r/r_{\text{cor}})$—such that the forcing is strong inside the resonance and rapidly weakens outside).

This gives us confidence that we can adopt the scalings above and robustly assume that there is indeed a characteristic radius (depending in detail upon orbital parameters) interior to which the gas will lose its angular momentum. This resonant structure of the angular momentum loss is actually quite convenient from an analytical perspective, as it means that more subtle issues of, e.g., the thermal pressure and state of the ISM, stellar and AGN feedback, and the exact mix of, e.g., gas and stars or density structure of the gas will not contribute significantly to determining which gas can or cannot lose angular momentum. Unlike, e.g., a self-generating bar in an unstable disk, there is no issue of stability analysis—the torques inside this critical radius (and the inducing perturbation) are sufficiently strong such that all the material therein loses angular momentum in a very short time (much less than a single orbital time, in practice).

For example, it is well known that in isolated cases, a pure gas disk is more unstable to gravitational perturbation than a stellar disk (see Christodoulou et al. 1995a, 1995b; Mayer & Wadsley 2004), however in the driven case this is not applicable: the distortion in the local stellar/gas distribution is caused by the secondary, not by, e.g., orbital “pileup” or instability in the primary. The location of the resonant radius is not determined by the internal structure of the primary (unlike in an isolated case, where it is determined by how, e.g., those orbits can overlap and where various stability criteria are satisfied), but rather by the orbital motion of the secondary (relative to the internal motion of the primary), and therefore knows nothing about, e.g., the gas to stars ratio, phase structure, and feedback situation in the primary. Inside this radius, the distortion is sufficiently strong that it does not matter whether one configuration or another is more or less prone to gravitational instability—the driving force (and therefore angular momentum loss) is large in any case.

Exactly what the pressure support of the gas inside this radius is may effect, e.g., how far it free-falls after losing angular momentum before shocking and forming a central starburst, but it will not change the fact that the angular momentum loss is efficient. Quantitatively, the torque is $\gg j_{\text{disk}} \Omega_d$ (as it must be in order for the gas to lose its angular momentum in much less than an orbital period); but, e.g., the pressure gradients resisting gas collapse cannot be larger than (in energetic terms) $\tilde{H} M_j V^2_\text{a} \sim \tilde{H} j_{\text{disk}} \Omega_d \ll j_{\text{disk}} \Omega_d$ (or else the disk could not be thin)—therefore whether or not there is even considerable pressure support or, e.g., thermal feedback or a modified ISM...
equation of state makes a negligible correction to the behavior seen in the simulations.

Before moving on, we would like to translate the general scaling above in terms of $\Omega_a$ and $\Omega_d$ into more convenient parameters. As noted above, $\Omega_a \sim v_{peri}/b$. We expect $b \sim 1-3 R_d$ for common parabolic cosmological orbits—as we discuss below, for orbits with larger $b$ that will eventually merge, all that matters in terms of the end product is the impact parameter of the final passage or two when the most dramatic forcing occurs, so even for initially larger passages, angular momentum transfer to the halo will ensure a value in this range toward the final stages of the merger.

Assuming a parabolic orbit, $v_{peri}$ will be given by the infall velocity from infinity, $\sqrt{(GM(1+\mu)/b)}$ (where $\mu$ is the merger mass ratio, discussed below). Because the merging systems are extended as $b \rightarrow 0$ these expressions should be replaced by a more complicated function of $b/R_d$ (for the case $b = 0$, the infall velocity asymptotes to the escape velocity from the center of the primary $\sim \sqrt{GM/R_d}$, which requires a numerical solution for an arbitrary density profile. In practice we find that we can interpolate between the limits $b = 0$ and $b \gg R_d$ quite accurately by replacing $b$ with $\sqrt{b^2 + R_d^2}$ (which also happens to be an exact solution for, e.g., a Plummer sphere density profile). Combining these factors, we find that (for the regime of typical interest)

$$\frac{\Omega_o}{\Omega_d} \sim \frac{v_{peri}}{v_c} \frac{R_d}{b} \cos(\theta)$$

$$= \sqrt{2} \frac{(1+\mu)}{[1+(b/R_d)^2]^{3/4}} \cos(\theta)$$

$$\approx 0.6 \cos(\theta), \quad (25)$$

where the last term comes from inserting a typical major merger mass ratio and $b \sim 2 R_d$. The orbital dependence is then—as we would expect—largely a function of the inclination angle $\theta$. Prograde orbits induce a strong bar response—despite the fact that in these mergers the orbital angular momenta are all aligned, we actually expect the most angular momentum loss and least efficient disk formation. Retrograde and polar mergers, on the other hand, despite having completely unaligned or canceling total angular momentum, should most efficiently form disks.

Inserting this dependence on orbital parameters into our previous derivation in Equation (17) allows us to effectively replace the part of $\Psi_{bar}$ which parameterized our ignorance of orbital parameters ($F(\ldots)$ in Equation (19)), giving

$$\Psi_{bar}(\theta, \hat{H}, \ldots) \propto \frac{1}{1 - \frac{\Omega_o}{\Omega_d} \cos \theta}. \quad (26)$$

In short, our previous derivation applies, but the orbital dependence is now explicit in $\Psi_{bar}$.

Revisiting Figures 7 and 8, recall that we included this orbital dependence in the predicted curves therein. For each orbit, the predicted curve is given by the solution for the gas mass within the critical $R_{gas}/R_d$ (Equation (17)) with the dependence on orbital parameters as in Equation (26)—we insert the appropriate orbital inclination $\theta$ and impact parameter $b$ for the two disks in the orbit and sum their expected $f_{burst}$ or $f_{disk}$ to derive the model prediction. The difference between the different orbits does not appear dramatic in Figure 7—but this is because the burst fraction $f_{burst}$ is plotted on a linear scale, suppressing the dependence on $\theta$ at small $f_{gas}$ (most of the visible dynamic range in the plot is at large $f_{gas}$—in this regime, however, the stellar bar is weak in any case because there is not much stellar mass in the disk—so the result is that much of the gas survives and becomes part of the disk, regardless of orbital parameters).

However, the difference between different orbits is much clearer in Figure 8, displayed on a logarithmic scale. At low $f_{gas}$, there is much more stellar mass in the disk than gas mass, so in principle the stellar bar could (if maximal) easily torque away all the angular momentum of the gas. Here, however, the orbital parameters become important in determining just how efficient this process should actually be. For the orbits close to retrograde ($\cos \theta \approx -1$), the scaling we have just derived suggests that $\Psi_{bar}$ should be suppressed by a factor $\sim 2$. But for orbits close to coplanar prograde ($\cos \theta \approx 1$), $\Psi_{bar}$ is enhanced by a factor $\sim 2$ to $3$—in other words, because the orbit is nearly resonant, the effective corotation resonance (the orbit interior to which the gas can efficiently lose angular momentum to the induced stellar bar) is moved out by a substantial factor, including a larger fraction of the disk gas (in those extremes, only the gas at very large radii survives the merger).

Again, Figures 7 and 8 demonstrate that the simple scalings based on our model provide an accurate description of the behavior in the full numerical experiments. Figure 9 combines these into a single plot—we compare the disk fractions in our simulations to the full expectation based on our derivation thus far as a function of both gas fraction and orbital parameters. As noted above, the agreement is good, with a reasonably small scatter.

These results are for four representative orbits spanning a reasonable range in orbital parameters—those for which we have a large number of simulations covering a wide range in the space of other parameters. We consider them first because this allows us to robustly determine that the predicted orbital scalings do not depend on, e.g., stellar mass, halo properties, feedback prescriptions, or other varied physics in the simulations. Having done so, we consider a more limited sampling of a much broader range in orbits given by Table 1 in order to survey the full dynamic range of orbital parameters.

Figure 10 shows the results of this. For a given suite of simulations with some particular orbital parameters, we first construct the correlation $f_{disk}(f_{gas})$ as in Figure 8. Rather than adopt some a priori model for the orbital dependence, we then fit the points in that correlation to a function of the form in Equation (19)—i.e., effectively fit for the normalization or “efficiency” of angular momentum removal, which we define as $(\Psi_{eff}(orbit))$. We compare this, for our ensemble of orbits, to our analytic expectation from the simple scaling in Equation (26) and to a full numerical solution (technically, for 1:1 mergers, we want to solve this separately for each disk and add the two, although just considering the primary is a good approximation for less major mergers). The agreement with our analytic model is quite good across the entire range of orbital parameters, implying that we have captured the most important physics of resonant interactions in this simple scaling.

We also compare this effective efficiency of disk destruction with the net specific angular momentum of the merger remnant (assuming pure addition/cancellation of the initial baryonic angular momenta of the disks and the orbital angular momentum).

The result is actually an anticorrelation: systems with aligned angular momentum vectors, e.g., coplanar prograde mergers being the extreme case, induce the most efficient bars and remove angular momentum most efficiently from the gas. Systems where the angular momentum vectors are misaligned (e.g., polar orbits)
and minor mergers with more limited studies of orbits (c, i, m, d), the red points are a study of major most well sampled orbits (e, h, k, f), shown in Figure 8, the purple points are <

we are interested in a comparison of orbital parameters here and not, e.g., gas The black solid line is the simple linear scaling ensemble of orbits, we expect values similar to those between our typical e and f simple manner predicted in Equation (26). Over a typical random cosmological constraint on \( \Psi \) place. This clearly emphasizes that it is not, in fact, any direct or antialigned (retrograde) actually leave the largest disks in

or antialigned (retrograde) actually leave the largest disks in place. This clearly emphasizes that it is not, in fact, any direct addition/cancellation of angular momentum that determines or enables disks to form in and survive mergers. Rather, the cases with the largest net angular momentum are most resonant, inducing the strongest resonant asymmetries in the merging pair, which most effectively drains angular momentum from the gas and leaves a compact, bulge-dominated remnant.

4.3.3. Dependence on Mass Ratio

The major remaining parameter to study is the merger mass ratio. Thus far, we have restricted our attention to equal mass 1:1 mergers, which allowed us to make several convenient simplifying assumptions. Nevertheless, most of our previous derivation applies. None of the scalings that we have explicitly derived up to now are dependent upon mass ratio. However, we have quantified the strength of the induced stellar and gas bars with the parameter \( \Psi_{\text{bar}} \), which we expect should scale with mass ratio. Moreover, we have made the assumption that the premerger disk stars are entirely violently relaxed by the merger. While this is a good assumption for 1:1 mergers, it is not true for minor mergers (a 1:10 mass-ratio merger, even with no gas, will clearly not transform the entire primary stellar disk into bulge).

First, consider this stellar component: there are a number of ways to derive the disturbance of the stellar component in the merger. The simple expectation is that the mass in galaxy \( M_2 \) which can be violently relaxed by collision with galaxy \( M_2 \) is proportional to \( M_2/M_1 \)—the net energy deposit, tidal forces, and the mass fraction brought in from a potentially disrupted satellite all scale in this manner. For simplicity, consider the case where the secondary \( M_2 \) is much smaller and more dense than the primary, and falls in on a nearly radial orbit in the final encounter (which is a good approximation, given the efficiency of angular momentum transfer from the orbit to the halo). Since we are assuming \( M_2 \ll M_1 \), treat \( M_2 \) as a point mass, and consider its final orbital decay, where it oscillates with rapidly decaying amplitude through the center of the primary with initial impact velocity \( v_i \approx v_\text{d} \) and damping spatial amplitude \( \ell_\text{max} \ll R_\text{d} \). At some instant, then, the secondary is at location \( (R', \phi', z') = (\ell \cos \theta, 0, \sin \theta) \) (we rotate such that the secondary orbit defines \( \phi = 0 \) without loss of generality).

A star in the primary disk at \( (R, \phi, z = 0) \) then feels some potential from the secondary (\( \equiv \Phi_\text{s} \)) and experiences a vertical deflection out of the disk \( \delta \Phi_\text{z} = (G M_2/R_\text{d}) \ell \sin \theta f(\ell/R) \), where \( f(u) = [1 + u^2 - 2u \cos \theta \cos \phi]^{3/2} \approx 1 \). We only interested in the time the secondary spends at \( \ell \sim R/v_i \) (when its much closer to the disk or further away, the vertical perturbation is weak), so it effectively acts for a time \( \delta t \sim R/v_i \) as it passes through \( \ell \sim R \) in its ringing about the center. If we know the full potential, we can solve for the deflection as a function of time and calculate the full acceleration of the disk stars at \( (R, \phi) \), which yields an effective net velocity deflection while the secondary is on one side of the galaxy of \( \delta v = (G M_2/v_i) (\ell_{\text{max}} \sin \theta/R) f(\ell_{\text{max}}/R) \). Deflection occurs when \( \delta v \approx v_i \) or larger, so if \( v_i = \tilde{v}(r) v_i \) (where \( \tilde{v} \) depends weakly on \( r \)) we substitute for \( v_c \equiv \sqrt{G M_1/R_\text{d}} \) here and in \( v_i \) we obtain the criterion \( G M_2/R_{\text{d}} \gtrsim v_c, v_\text{i} \approx \tilde{v}_\text{i} \approx G M_1/R_\text{d} \), i.e., (rearranging) \( R/R_\text{d} \gtrsim M_2/M_1 \sin \theta = \mu \sin \theta \). The \( \sin \theta \) dependence comes because we considered only vertical deflection of stars (i.e., some heating to \( v_\text{z}^2 \))—a coplanar orbit (in this limit) will obviously induce no such heating, but will introduce deflections in the radial direction (heating \( v_\text{R}^2 \)). We can repeat our derivation considering where these deflections are significant, and find (as one would expect) \( R/R_\text{d} \gtrsim \mu \cos \theta \).

So, the absolute mass fraction scattered should be more or less angle-dependent, although the orbital anisotropy \( \beta_\parallel \equiv 1-v_\text{R}^2/v_\text{c}^2 \) will depend significantly on the orbital inclination \( \theta \).
For the case of a thin Mestel (1963) disk with no bulge, we can solve these equations exactly and obtain the simple solution that a merger with a secondary $M_2$ scatters exactly $M_2$ worth of stars in the primary, completely independent of the inclination $\theta$ (but with an anisotropy $b_\ell(\theta) \sim 1 - 2 \frac{\sin^2 \theta}{(1+2 \cos^2 \theta)^2}$, although this ignores a proper treatment of further mixing as the perturbed stars interact with each other, and thus does not reproduce orbits quite as radial as seen in simulations). The full numerical solutions for arbitrary cases yield the general result that, when inside a radius that encloses a mass $\sim M_2$ in the primary, then the presence of the mass $M_2$ is a significant perturbation, which scatters those stars in the primary—i.e., deflections occur rapidly, so the stars violently relax. At larger radii, where $M_{\text{enc}} \to M_1 > M_2$, the motion of the $M_2$ secondary at the center is a small perturbation. The disk at these radii is perturbed adiabatically by the motion of the secondary, which can induce some warps and/or disk heating, but will not violently relax the stars. Reversing this derivation for the secondary, it is trivial that essentially all the mass in the secondary (we ignore stripping of the tightly bound stellar mass) will be violently relaxed. So the total mass violently relaxed will be $\sim M_2$ (in the primary) plus $M_2$ (the secondary), out of a total mass $M_1 + M_2$—i.e., in terms of the mass ratio $\mu \equiv M_2/M_1$, the fraction of the premerger stellar disk mass which is destroyed and turned into bulge is

$$f_{\text{rel, disk}} \approx \frac{2 \mu}{1 + \mu}. \quad (27)$$

Technically this assumes the systems are initially pure disk, but the corrections if they have pre-existing bulges are not large (generally smaller than the simulation-to-simulation variation; although we discuss them in more detail in Section 5), so this is a reasonable approximation for general cases.

Now, consider the gas. It turns out that a similar linear scaling in Equation (27) is found for how the gas mass in the starburst (i.e., the fraction of the premerger stellar disk mass which is destroyed and turned into bulge) scales with $\mu$. The details we are interested in is dominated by modes with $k \sim 1/R_d$; but we can, for example, treat $\Phi_g(k R)$ as the potential generated by a point source (appropriate for, e.g., the small mass-ratio limit) at the impact parameter $b$ and expand the wave modes appropriately, and integrate over the modes in the disk to determine the bar strength (i.e., the total mass effectively contributing to the bar). In any case, up to a numerical constant that is weakly sensitive to the mode structure, we obtain

$$f_{\text{rel, gas}} \approx \frac{2 \mu}{1 + \mu}. \quad (28)$$

where again $b$ is the impact parameter and the $1 + [b/R_d]^3$ term effectively allows for the interpolation between the case of an orbiting point mass and a penetrating encounter (see Binney & Tremaine 1987).

As noted in Section 4.3.2, if we are just interested in the end product of a merger—i.e., we do not care what happens on each passage separately as the companions lose angular momentum, but only in the surviving disk fraction and total burst fraction—then we are not interested in some initial impact parameter $b$ but only in the impact parameter on the final passages close to coalescence, when angular momentum loss has made the orbits nearly radial and the forcing is strong. We can see directly from Equation (29) that the forcing is dramatically suppressed by a factor $\sim (b/R_d)^3$ on earlier, large-impact parameter passages, so these can be effectively ignored in calculating the remnant properties (we confirm this is true in a sample of simulations with much larger $R_{\text{pert}}$). Eventually, for any systems which are destined to merge, angular momentum transfer yields a nearly radial orbit with $b \sim R_d$ and this is where most of the forcing occurs, so the remnant solution is effectively given by ignoring the $b$ dependence above (technically summing over each passage with the appropriate $b$ is possible, but in practice we obtain the same result to within the simulation-to-simulation scatter by assuming $b \to 0$ in Equation (29)). The final dependence on any initial impact parameter $b$ is therefore weak, so long as the systems are bound to merge. The dependence on mass ratio $\mu$, however, is fixed.

If the mass enclosed is linear in $\Psi_{\text{bar}}$ (the case for, e.g., the Mestel (1963) disk and an isothermal sphere, and not a bad approximation for the regime of interest for an exponential disk), we then have a similar result for the gas as the stellar distribution: the secondary induces a burst of mass $\propto M_2$ in the primary $M_1$. Reversing the derivation, the primary (since it is larger, so $M_1/M_2 > 1$ induces a burst (assuming the two have similar initial gas fractions) $\propto M_2$ in the secondary (i.e., bursting all its gas). The net burst mass $\propto 2M_2$ relative to the remnant mass $M_1 + M_2$ is then

$$f_{\text{burst}} \propto \frac{2 \mu}{1 + \mu}. \quad (30)$$

This is generally applicable for mergers; however we will note below that, because they do not coalesce (and therefore do not
Figure 11 tests this prediction that simulations spanning a range in mass ratio from $\mu = 0.1$–1. For a given set of orbital parameters (fixed), we plot the disk and burst fractions ($f_{\text{disk}}$ and $f_{\text{burst}}$) of the remnant, as a function of the immediate premerger gas fraction $f_{\text{gas}}$, as in Figures 7 and 8. For the 1:1 mergers, we plot our expectation based on the simple scaling in Equations (17)–(26), including the dependence on $\theta$ and $\mu$ from Equation (27)—so some (considerable) fraction of the gas that loses its angular momentum (fraction of $f_{\text{gas}}$) and fraction of the premerger primary stellar disk turned into bulge (fraction of $(1 - f_{\text{gas}})$) are suppressed by a factor $\sim \mu$.

For each of the orbits surveyed (and the range in, e.g., absolute masses, gas fractions, and feedback prescriptions in our minor merger simulations), this simple rescaling according to the merger mass ratio provides a good approximation to the behavior in the full hydrodynamic experiments. Both the total surviving disk fraction (which reflects both the ability of the premerger stellar disks and the premerger gas to survive the merger) and the burst fractions (which reflect only how much of the gas survives/loses angular momentum) are accurately predicted, suggesting that our derivations are reasonable for both the dissipational and dissipationless components of the galaxy.

Figure 12 summarizes these results. We first compare the final disk fraction in the simulations to our prediction including, e.g., the dependence on gas content and orbital parameters but without any accounting for mass ratio (assuming all mergers are just as efficient as a 1:1 merger). Unsurprisingly, this works for the 1:1 mergers, but is a terrible approximation to mergers of very different mass ratios. We then compare allowing for the same scalings but including the predicted mass ratio dependence. The agreement between full simulation and our simple analytic expectations is good—with a scatter for the high disk fractions typical of intermediate and minor merger remnants as low as $\sim 20\%$.

One important caveat here is that, for mergers of increasingly small mass ratio $\mu$, the merger timescales become long. At the smallest mass ratios we consider, $\sim 1:10$, this timescale may become sufficiently long that the secular (i.e., self-amplifying) instability/response of the disk may become important over the duration of the merger. It is not entirely clear what the response of such an (initially driven) system will be; whether or not, for example, the driven nonaxisymmetric modes will remain locked to their driver (the secondary orbit) or decouple and move at the pattern speed dictated by the internal stability properties of the disk. This competition between secular processes (more sensitive to, e.g., the detailed structure, rotation, and pressure support of the disk) and merger-driving in this regime probably contributes to some of the increased scatter in burst fractions seen in Figure 11 at the lowest $\mu$. For this reason, it is reasonable to restrict a definition of “mergers” to this mass ratio and more major interactions: at smaller mass ratios, secular/internal processes (even if initially driven by interactions) may be more important than the direct driving from the interactions themselves (or at least operate on comparable timescales).
characteristic scale length of the system. But in detail, the appropriate “impact parameter” is really the actual distance of closest approach, which is usually somewhat smaller than the distance of approach estimated from infinity (the formal impact parameter definition), as some angular momentum is already lost. Moreover, in detail, is the appropriate $R_d$ the exponential scale length? The half-mass radius? Any such radii are of course closely related, and all of these uncertainties in the exact definition change the term $b/R_d$ only at the factor $\sim 2$ level, but since the dependence $\sim (b/R_d)^3$ is fairly strong, this is important on a quantitative level for these BH situations.

In practice, we find that using the impact parameter $b$ defined as the halos approach (i.e., neglecting detailed resonant loss of angular momentum) and taking $R_d$ to be the half-mass radius of the system works well in a mean sense. The results of this exercise are shown in Figure 13. We plot the fraction of the gas available at the time of a first passage or BH encounter which is consumed in the induced burst (we define the strength of the induced burst by integrating the star formation excess over the interpolation between the pre- and post-flyby SF rates; see Cox et al. 2008, for details), as a function of the gas content, for different orbits as in Figure 8. We predict that the efficiency of channeling gas into the burst should scale as $\sim (1 - f_{\text{gas}})$, as before, and that the scaling with orbital parameters should be similar. We also show, for cases with otherwise identical gas fraction at the time of first passage and the same orbits as Figure 11, how this scales with merger mass ratio (again, expected to be the same as that we derived above). Altogether, adopting our previous estimates, but renormalizing appropriately for the impact parameter of the passage (here $b/R_d \approx 1$, as we defined it above) yields a good approximation to the typical behavior in our simulations. We also test the behavior as a function of impact parameter, and find that our simple scaling is a reasonable approximation, yielding rapidly diminishing bursts as the impact parameter is increased to $b \gg R_d$ (at some point here, our estimates from the simulations become ambiguous, as a $\sim 1\%$ enhancement in star formation is below the level of random fluctuations in isolated disks). We have also checked whether or not the pre-flyby stellar disks are destroyed—as expected, they are left more or less intact by BH encounters. The disks may be heated, and in fact some “pseudobulge” can form from the buckling of the bar induced in the stars, but we are not attempting to predict or study pseudobulge formation here (rather considering it, as is often the case in observations, to be fundamentally still part of the stellar disk rather than part of a violently relaxed “classical” bulge). In an average sense, then, our derived scaling is generally applicable.

However, the details of exactly how the approach proceeds will introduce considerable scatter in the amount of burst triggered on first passages and in BH encounters. This is plain in the large (factor $\sim$ a few) scatter in Figure 13. Further, details such as the structure of the bulge are increasingly important in the limit of weak interactions, where distortions in the potential of the primary that would trigger gas inflows can be suppressed by the presence of a larger bulge (and note the caveat from Section 4.3.3, that the secular/internal response of the disk will become relatively more important in weaker interactions with smaller mass ratios and larger impact parameters). We therefore expect in general that our predictions can be quite broadly applied, but are less robust for any specific case if it is a single BH as opposed to an integration over a full merger. Fortunately, in the case of systems that will actually merge,
these details tend to average out or be unimportant, yielding the relatively small scatter we have seen in our previous predictions. In those cases, we do not need to be too concerned with the exact details of the impact approach, nor the structural details of the galaxy (in particular because our predictions are for integral quantities at the end of a merger, various effects will tend to cancel out, for example, retaining more gas on first passage will yield a larger supply for the second burst, etc.).

4.3.5. Independence from “Feedback” Physics

Our derivation of the torques causing gas to lose angular momentum in mergers is purely dynamical. All else being equal (i.e., for systems with the same gas content and dynamical structure at the time of the final merger), we therefore expect that the detailed physics of, e.g., “feedback” from supernovae, stellar winds, and AGN activity should make little difference.

Figure 14 demonstrates that this is indeed the case. We compare the starburst and surviving disk gas fractions of merger remnants, relative to those predicted by our simple dynamical model as a function of the merger mass ratio, orbital parameters, and gas content at the time of the merger, for suites of simulations with two different prescriptions for supernovae feedback and the effective equation of state of the ISM. In terms of our $q_{\text{EOS}}$ parameter (see Section 2), we compare $q_{\text{EOS}} = 0.25$ simulations (a nearly isothermal equation of state with effective temperature $\sim 10^{4}$ K) to $q_{\text{EOS}} = 1$ simulations (the “full” stiff Springel & Hernquist (2003) equation of state, with effective temperature $\gtrsim 10^{5}$ K at the densities of interest here). There is no significant systematic offset between either the median result or the scatter about our simple analytic expectation. At most, there may be a $\sim 20\%$ systematic offset, in the sense that more highly pressurized systems ($q_{\text{EOS}} = 1$) have slightly more gas survive—a small offset like this is expected because the bars in these cases are slightly more “puffy” (essentially the same as a slightly thicker disk—for which we derive an analytic expectation in Equation (20) that yields an expected $\sim 10\%$–$20\%$ difference at most based on the full possible range of $q_{\text{EOS}}$). In any case, such an offset is small relative to other systematic uncertainties in disk structure and the scatter about the median predictions.

In Figure 15, we perform a similar exercise for cases with and without central supermassive BHs (we have also examined initial BHs with varying initial masses from $\lesssim 10^{5} M_{\odot}$ to...
with a mass-loading efficiency to the stellar feedback implicit in our multiphase ISM model) of starburst-driven winds where winds are launched (in addition $10^7 M_\odot$), and cases with or without a simple implementation of starburst-driven winds where winds are launched (in addition to the stellar feedback implicit in our multiphase ISM model) with a mass-loading efficiency $M_{\text{wind}} = \eta_w M_\ast$ relative to the SFR $M_\ast$ and energy loading efficiency $\epsilon_w$ relative to the total energy (for a Salpeter 1955 IMF) available for supernovae (sampling the range $\eta_w \sim 0.01$–10 and $\epsilon_w \sim 0.1$–1; see Cox et al. 2008). Shown in Figure 15 are our fiducial weak winds ($\eta_w \sim 0.01$, $\epsilon_w \sim 0.0025$), cases with moderate mass loading into very fast winds ($\eta_w = 0.5$, $\epsilon_w = 0.25$, yielding a wind launch speed $\sim 800$ km s$^{-1}$), and cases with high mass loading but correspondingly slower wind velocities ($\eta_w = 2.0$, $\epsilon_w = 0.0625$, yielding a wind launch speed $\sim 200$ km s$^{-1}$). In all these cases we find a similar result: at otherwise fixed properties at the time of merger, feedback makes no difference to our conclusions.

The reasons for this are described in Section 4.3.2 and in more detail below (Section 4.3.7). Recall, the distortion in the primary is driven by the secondary and as such depends only on the gravitational physics of the merger. Given this distortion, the gravitational torques within some characteristic radius are sufficiently strong to remove the angular momentum from the gas in much less than an orbital time. Feedback, then, insofar as it changes the effective pressurization or equation of state of the gas or drives a wind, is largely irrelevant: because the angular momentum is removed in a timescale much shorter than the orbital time, the gas (regardless of the strength of feedback) cannot dynamically respond with these hydrodynamic forces, but must essentially free-fall into the center of the galaxy where the starburst is triggered. The radius interior to which the torques are strong is not a function of, e.g., the stability of the galaxy to perturbation, because it is not an instability in the first place, but a driven distortion in the system. Moreover, the entire system is strongly in the nonlinear regime for the time of interest (when we consider interactions of the magnitude simulated: mass ratios $\sim 1:8$ and more major mergers)—no amount of making the system more robust against linear instability would be sufficient to avoid a strong gravitational distortion in the violent coalescence surrounding the actual merger (this must be so, because the distortion occurs where the disturbances in the

Figure 14. Effects of feedback on disk survival in mergers. We show the distribution in $f_{\text{burst}}/f_{\text{burst, pred}}$, i.e., the burst mass fraction, relative to that predicted (or equivalently, the mean in our simulations) for the given premerger gas fraction and orbital parameters, for simulations with two different ISM gas feedback prescriptions (different effective equations of state $\eta_{\text{gas}}$). We also show the corresponding (but measured differently) disk mass fractions $f_{\text{disk}}/f_{\text{disk, pred}}$. There is perhaps a small offset in the sense expected (a stiffer, higher feedback equation of state for the ISM suppresses bursts by an average factor of $\sim 1.1–1.2$), but this is much smaller than the simulation-to-simulation scatter. For a given gas content at the time of the merger, then, feedback makes almost no difference (true for AGN feedback and starburst winds as well).

(A color version of this figure is available in the online journal.)

Figure 15. As Figure 14, but comparing the distribution of disk fractions in merger remnants relative to our simple predictions as a function of stellar wind and quasar feedback prescriptions. We plot the distribution of disk fraction $f_{\text{disk}}$ in simulations relative to our predicted $f_{\text{disk}}(f_{\text{gas}}, \mu, \theta)$ (i.e., our calculation as a function of immediate premerger gas fraction, orbital parameters, and merger mass ratio). Left: varying starburst-driven wind prescriptions. We compare our usual weak stellar wind launch speed $\sim 800$ km s$^{-1}$ and launch velocity of $\sim 200$ km s$^{-1}$). Right: simulations with and without feedback from accreting BHs. For otherwise fixed merger parameters (orbit, mass ratio) and disk properties (cold gas content, mass profiles), these feedback prescriptions make no different to the starburst or surviving disk gas fractions. They will, however, change the gas content, consumption, and distribution leading into the merger.

(A color version of this figure is available in the online journal.)
potential are greater than order unity—for hydrodynamic forces to resist distortion, they would have to be stronger than large-scale gravitational forces in the equilibrium system, negating the concept of a rotationally supported thin disk). So what matters is instead where that coalescence occurs and how long it introduces such a strong distortion, relative to, e.g., the local dynamical or orbital time of some disk element, giving rise to the simple dynamical criteria for angular momentum loss developed here.

That is not to say that for fixed initial conditions (significantly premerger or, e.g., at first passage), feedback will not change the result. There are two primary means by which feedback can indirectly have a strong influence on disk survival:

1. Retaining gas (lowering the stellar mass fraction): as has been demonstrated in a number of works (Weil et al. 1998; Sommer-Larsen et al. 1999, 2003; Thacker & Couchman 2000, 2001; Governato et al. 2007; Robertson et al. 2006a; Springel et al. 2005b; Springel & Hernquist 2005; Okamoto et al. 2005; Scannapieco et al. 2008), these forms of feedback can have dramatic implications, in even a short time period, for the rates at which cooling of new cold gas from the halo and consumption of existing gas by star formation proceed. In cases with no feedback, star formation may exhaust gas efficiently, leading to predicted systems that are much more gas-poor at the interesting time of the final merger—according to our model, then, these will not be able to form disks as efficiently as more gas-rich systems. In cases with strong feedback from, e.g., star formation to lower the effective star formation efficiency and recycle gas, the predicted gas fractions at the time of merger (from some gas-rich initial conditions) could be much higher. Inclusion of stellar and supernova feedback responsible for injecting energy and turbulent pressure into the ISM may also be necessary to prevent the onset of clumping and disk fragmentation in isolated gas-rich cases, enabling the stable existence and evolution of quiescent gas-rich disks (see Springel & Hernquist 2003; Robertson et al. 2004). In short, feedback may be critical to give rise to high gas fractions in the first place, which we have shown have dramatic implications for the survival of disks—but for a given gas fraction (however that comes about in the first place), the results of the merger will be (in the short term) independent of feedback.

2. Changing the spatial distribution of gas (“kicking gas out” of $R_{\text{max}}$): recall, our derivations demonstrate that it is not necessarily a fixed fraction of gas that loses its angular momentum; rather (see Equation (17) and Sections 4.3.1 and 4.3.2) it is the mass inside some radius $R_{\text{max}}/R_d$ relative to that of the stellar disk (characteristic radius $R_d$) which will lose its angular momentum. If some form of feedback can change the spatial gas distribution, then, it could have dramatic implications for disk survival. We have used the radius $R_{\text{max}}$ to estimate the mass fraction that will burst by assuming that the gas density profile is broadly similar to that of the stars (which is true in our simulations, given their feedback prescriptions). But one could easily imagine the extreme limit, where some strong feedback keeps all the gas at large radii $r \gg R_{\text{max}} \sim R_d$ (i.e., a case in which there is a large hole in the gas distribution, or in which the gas is at least much more extended than the stellar distribution)—the stellar disk torques only act effectively within $R_{\text{max}}$, so only a tiny fraction of the gas in such a case would lose its angular momentum. Especially at high redshift, this may be important in avoiding overcooling and the formation of too much bulge mass in many systems (see Robertson et al. 2004; Governato et al. 2007; D’Onghia et al. 2006; Ceverino & Klypin 2007; Zavala et al. 2008). Again, we stress that for a given gas density profile at the time of merger, our calculations are independent of feedback; but if feedback alters the gas profile—keeping the gas at radii $r \gg R_{\text{max}}$, then it will largely survive the merger.

### 4.3.6. Exceptions and Pathological Cases

We have derived a general model for how disks are destroyed in mergers and shown that it applies to a wide range of gas fractions, orbital parameters, galaxy mass ratios, and prescriptions for feedback and gas physics. However, there are some pathological cases of more than academic interest, as these can explain some small differences with previous results as well as illustrate the important physics in our model.

For example, consider the starburst mass fraction and surviving disks in our “h” orbits: i.e., a prograde–prograde, coplanar merger of two disks. In this case, the angular momentum vectors of both disks and the orbital angular momentum are all perfectly aligned. Naively, one might then expect that these unique cases would create the largest disks. In fact, the opposite is true. This is largely for the reasons we outline in Section 4.3.2—the alignment of angular momentum vectors means that the system is in near-perfect resonance, so it excites the largest tidal and bar asymmetries that rapidly drain the gas of all angular momentum. As we have shown, the much larger space of less-aligned orbits is in fact more favorable to disk survival.

However, while the perfect resonance means that the bar efficiency $\Psi_{\text{bar}}$ is large, the amount of mass in the stellar bar still scales as $1 - f_{\text{gas}}$, so the burst fraction should vary as $f_{\text{gas}}(1 - f_{\text{gas}})$ in our simple model (i.e., we would still expect that a 100% gas disk would have no stellar bar, hence no burst, as we have seen for more representative orbits in Figure 7). In fact, though, we typically find in these cases that the burst fraction seems to scale as $f_{\text{burst}} = f_{\text{gas}}$ all the way to high values of $f_{\text{gas}}$—in short, almost all the gas always bursts—there is no suppression by a $1 - f_{\text{gas}}$ factor as would be expected if the stellar bar were doing the torquing.

The reason for this is simple—again, the orbits here are perfectly coplanar and in resonance; so this is the one case where the secondary galaxy as a whole can directly act as an efficient torque on the gas. In short, because the systems are perfectly coplanar and in a resonant orbit, the entire secondary galaxy (all baryons and dark matter within the stellar $R_c$) acts directly to introduce a nonaxisymmetric potential perturbation (the secondary itself plays the role of the bar). So because of this, to an even greater extreme than our scalings for more general orbits would predict, this narrow range of orbits is pathological and biased against disk formation. However, understanding why this is the case, we can check and explicitly show that it is not so for more general orbits, even nearly prograde–prograde orbits (such as case e)—in all those cases, even those just slightly out of coplanar resonance, the stellar bar is indeed the primary source of torque, and our assumptions are justified. This example therefore nicely illustrates what the consequences would be if our fundamental assumptions were not true, as well as showing why they are in fact true for nonpathological cases.

Another pathological case of interest is one in which the disks are 100% gas at the time of merger. Here, as we have said, our simple model predicts no starburst or angular momentum loss. In practice, there will still be some loss of angular momentum owing to direct cancellation in, e.g., shocks between the disks;
but as discussed in Section 4.2, there will also be the possibility of some gain owing to the angular momentum of the merger. In fact, over the range in mass ratios $\mu \sim 0.1$–1, for a range of typical impact parameters $b \sim 0.5$–5, the expected final specific angular momentum after cancellation is approximately equal to the initial specific angular momentum of the primary (with $\sim 20\%$ scatter). Cancellation is therefore inefficient. A random distribution of orbits might negate $\sim 20\%$ of the angular momentum in $\sim$ half the systems merging, but will leave $\sim 80\%$–$100\%$ of the disk intact. Even these cancellations, we find in detail, do not generally yield a starburst in the same manner as a merger-induced bar, but lead to moderate disk contraction (and an equal number of mergers will scatter toward the opposite sense leading to disk expansion, keeping a mean specific angular momentum that is constant). They do not cause a starburst because, if two random parcels or streams of gas shock and lose angular momentum, the alignment and relative momenta would have to be near-perfect for them to lose, say $95\%$ of the angular momentum and fall all the way to the central $\sim 100$ pc where a nuclear starburst would occur. Rather, they will lose some fraction of order unity of their angular momentum, fall in to a radius smaller by a factor $\sim 2$–3 (but not to very small radii), and continue to orbit. Without the bar that can continuously drain angular momentum, the true burst is indeed inefficient.

Although we show in Section 4.3.5 that the physics of interest are generally independent of feedback prescriptions, there are some pathological feedback regimes. These are discussed in detail in T. J. Cox et al. (2009, in preparation); here, we outline the pathological behavior. If, e.g., starburst-driven winds are implemented with extreme efficiencies $M_{\text{wind}} \gg M_\star$, and with moderate to large velocities $\gtrsim 200$ km s$^{-1}$, then there is no definable “starburst” in the simulations any more, even when the gas loses angular momentum—indeed, it becomes almost impossible to trigger starbursts by any mechanism. This is because the feedback is so extreme that any parcel of gas that begins forming stars above some threshold rate is immediately blown apart and drives away all the surrounding gas. However, observations suggest that these cases are almost certainly not relevant—observationally inferred mass-loading factors of winds are well below the predicted threshold were we see this behavior (see Veilleux et al. 2005; Martin 1999, 2006; Erb et al. 2006a; Sato et al. 2008), and moreover the ubiquity of starbursts and recent starburst remnants in observed gas-rich major mergers (e.g., Soifer et al. 1984a, 1984b; Scoville et al. 1986; Sargent et al. 1987, 1989) implies that feedback, while still potentially efficient, is not able to “self-terminate” a starburst before it even begins (this is in fact directly confirmed in observations of outflows in ongoing massive, merger-induced starbursts; see Martin 2005). A similar pathology can appear if we include extreme coupling of BH feedback to the galaxy gas (e.g., allowing $100\%$ of the BH accretion energy to couple efficiently), but this is also ruled out observationally, both for the arguments above (starbursts exist, and the winds seen are not so enormous; see the discussion in T. J. Cox et al. 2009, in preparation, Hopkins et al. 2005a, 2005b, 2005c, 2005d, 2006), and because such a prescription yields BH masses orders-of-magnitude discrepant from the observed (Ferrarese & Merritt 2000; Gebhardt et al. 2000) $M_{\text{BH}} - \sigma$ relation (see Hopkins et al. 2007a, 2007b).

### 4.3.7. Longer-Lived Perturbations: Relation to Secular Evolution

Thus far, we have focused on activity during the merger, roughly defined as the short timescale $\sim 10^8$ yr following first passage and coalescence. In this regime, we have shown that (for typical conditions) the dominant source of angular momentum loss is the torque on gas from stars in the same disk. However, it is well known from studies of isolated barred galaxies (e.g., Weinberg 1985; Hernquist & Weinberg 1992b; Friedli & Benz 1993; Friedli et al. 1994; Athanassoula 2002a; Athanassoula & Misiriotis 2002; Weinberg & Katz 2007; Kaufmann et al. 2007; Foyle et al. 2008) that a long-lived bar (regardless of whether the bar is purely stellar or purely gaseous) will exchange angular momentum with itself (or, e.g., gas/stars further out in the disk) and the dark matter halo, allowing for further angular momentum loss and building a central bulge or “pseudo”-bulge (Patsis & Athanassoula 2000; Athanassoula 2002b; Mayer & Wadsley 2004; Berentzen et al. 2007). Here, we discuss the relation of this process to what we have described in our merger simulations: in general, we find that it (while potentially very important for the long-term evolution of the disk and bulge masses and structure) is a second-order effect within the merger itself, and on longer timescales is more appropriately considered an independent, secular evolution process (despite being initially triggered by a merger), whose study is better described in simulations of idealized and long-lived bars.

As discussed in Section 4.3.2, there is a limit to how far the analogy to barred galaxies can be drawn. Recall, we use the term “bar” more generally to represent a quadrupole moment or nonaxisymmetric distortion in the stellar disk: it does not necessarily (and, especially after second passage, usually does not) morphologically resemble isolated barred spirals and may not even have an $m = 2$ mode structure. Critically, the distortion is driven externally by the gravitational perturbation of the secondary orbit—it is not the result of an instability within the primary. As we note in Section 4.3.2, this already gives rise to a couple of important distinctions: because the distortion is driven by the orbital motion of the secondary, it has a characteristic frequency (and the corresponding radius) internal to which angular momentum loss is very efficient (determined entirely by the gravitational properties and relative motions of the systems, not the subtleties of their internal orbital structure), regardless of properties of the primary (e.g., gas phase structure, feedback, etc.) that might otherwise make the system more or less stable to the development of internal instabilities.

It is straightforward to estimate the relative importance (over a short timescale after the merger) of angular momentum loss from the gas to the shared stellar bar/distortion induced by the merger, versus that to itself and the dark matter halo (the standard secular scenario, which other than the initial driving in the merger, will not be driven by the relative gravitational motions but by the more standard bar stability and spin-down criteria). Approximating the gas as a rigid, thin bar of mass $M_{\text{bar, gas}} \approx f_{\text{gas}} M_{\text{bar}}$ and radius $R_{\text{bar}} \sim R_\sigma$, we can estimate the specific torque from the remaining gas disk and halo in the dynamical friction limit, following Weinberg (1985; for more detailed solutions, which ultimately give similar results, see Hernquist & Weinberg 1992b; Athanassoula 2003; Weinberg & Katz 2007): $dj/dt = -4\pi \alpha G^2 M_{\text{bar, gas}} \rho(R_{\text{bar}}) v_{\text{bar}}^2$, where $\rho(R_{\text{bar}})$ is the background density and $\alpha \sim 1$ is a numerical constant (depending on the exact shape of the bar, potential, and phase-space distribution of the background). If the “background” is a Mestel (1963) disk or isothermal sphere, this becomes $-2\alpha G M_{\text{bar, gas}} / R_\sigma$. Compare this to our Equation (12) for the instantaneous torque on the gas bar from the stellar bar: $-G M_{\text{bar}} \sigma^2/(R_{\sigma} \sqrt{\sin^2 \phi_0 + H^2})$. Removing the common factors, the gas/halo torque goes as $\sim f_{\text{gas}}$, whereas that from the
stellar bar goes as \(\sim (1 - f_{\text{gas}})/\sqrt{\sin^2 \theta_0 + H^2} \sim (1 - f_{\text{gas}})/\phi_0\) (because \(\phi_0 \sim H \ll 1\)). In short, the torque from the gas disk and halo goes as \(f_{\text{gas}}\) because it is a second-order resonance effect (amplified and trading off with the gas bar), whereas the stellar bar strength goes as the stellar mass fraction \((1 - f_{\text{gas}})\), but boosted by a factor \(\sim 1/\phi_0\) representing the small angle of offset between the two bars—i.e., the stellar bar is in much closer spatial proximity (in particular in spatial alignment in the disk plane) to the gas bar.

This simple comparison gives a reasonable quantitative prediction of the relative torques exerted by the halo and stellar disk in our simulations. Essentially, we have just rederived the well-known fact that the timescale for a bar to damp its own angular momentum via resonant interactions with itself and/or the halo is some number \((\sim a)\) bar rotational periods (Athanassoula & Misiriotis 2002; Athanassoula 2003; Weinberg & Katz 2007; Kaufmann et al. 2007; each bar rotational period being \(\sim 1 - 2\) times the disk rotational period), whereas in the typical mergers the gas is drained of angular momentum by the much stronger local torques on a timescale much shorter than an orbital time, allowing it to more or less free-fall into the galactic center. Comparing these timescales gives a similar ratio of torque strengths. Obviously, as \(f_{\text{gas}} \rightarrow 1\), the torque from the halo must eventually dominate, but this will not happen until \(f_s = (1 - f_{\text{gas}}) \lesssim \phi_0 \sim 0.1\) (given typical bar lags of \(\sim a\) degrees or the ratio of the timescales above).

In practice, such a situation is somewhat contrived (it is very difficult to maintain a disk with a true \(\gtrsim 90\%\) gas fraction), and unlikely to be of broad cosmological relevance (we have no simulations in this regime with which to compare, in fact, because even initially 100\% gas disks with low star formation efficiencies will be \(\lesssim 80\%\) gas by the time of the actual merger). However, our Equation (7) can be trivially modified to include these effects: the \((1 - f_{\text{gas}})\) term should be replaced with a more appropriate \((1 - f_{\text{gas}} + \epsilon_h)\), where \(\epsilon_h \sim \phi_0 \sim 0.1\) represents the contribution of angular momentum loss to the halo and outer gas disk during the merger. This exchange of angular momentum, therefore, sets some minimum bulge mass (with mass fraction \(\sim 10\%\)) that would form even in a pure-gaseous disk merger.

By comparing the relative instantaneous amplitude of the torques from the halo/gas and stellar disk, we are comparing how important they each are in the loss of angular momentum from the gas over the same (relatively short) merger timescale. More important is the fact that, in a gas-rich case where the distortion to the stellar distribution may be an inefficient torque, the gas bar could be long-lived and continue to lose angular momentum over longer timescales. It is not necessarily clear that this would happen, however, a number of studies suggest that gas and stellar bars become self-damping once a central mass concentration (i.e., a nuclear starburst triggered by gas inflows, in this case) is in place with a mass fraction larger than a few percent (Bournaud & Combes 2002; Berentzen et al. 2003, 2004, 2007; Athanassoula et al. 2005; but see also Kaufmann et al. 2007). In such a case, we again arrive at the conclusion that these processes set a minimum bulge mass from a large bar-inducing perturbation, but do not dominate the creation of much larger bulges in mergers.

Regardless of this effect, it is not clear that a bar can survive a substantial merger: recall, the distortion following second passage and coalescence resembles a bar only in that it introduces a rotating quadrupole distortion in the disk potential (allowing us to describe it as a “bar” for analytic convenience), not necessarily in its structure or longevity (it does not necessarily share the orbital “pileup” that allows a bar to survive), and moreover the actual coalescence of the galactic nuclei will disturb any bar structure that may be present. Quantitatively, we find that our remnants rarely have significant long-lived \(m = 2\) modes in the stars or gas—the mode amplitude tends to damp after merger on a timescale \(\lesssim 10^8\) yr (i.e., the free-fall or dynamical time, much slower than the typical significant number of orbital times for standard bar self-braking). This is similar to the conclusions in the bar studies of, e.g., Bournaud & Combes (2002) and Berentzen et al. (2007), who find that the combination of the formation of a bulge/small central mass concentration from gas inflows and the disturbance/heating to the bar itself in interactions prevents even gas-rich systems from maintaining or very rapidly re-forming a bar after a significant merger (as opposed to a BH passage, which may more efficiently induce long-lived bars; see Berentzen et al. 2004). That is not to say a bar may not form in the re-formed remnant disk, but such a bar would arise in a standard secular fashion, and should be considered in the context of the long-term secular evolution of the merger remnant.

If, however, the potential distortion survives the merger to form a stable bar, it can certainly be important to the long-term evolution of the system and buildup of the bulge. However, this case is outside the scope of this paper, and should be more appropriately considered as subsequent evolution of the remnant (albeit with an initially merger-induced bar). This is because the timescale for the bar to lose angular momentum and contract is some number of rotational periods—so the gas losing angular momentum will slowly spiral inward in some number of orbital periods (turning into stars and possibly being ejected by feedback as it does so), rather than free-falling into a central burst in a time much less than an orbital period. The end result of such angular momentum loss can resemble a bulge (Mayer & Wadsley 2004; Debattista et al. 2004), although the expectation of rotational support and “disky-ness” in the material lead to it more likely being a “pseudobulge” typical of secular processes (Combes et al. 1990; Kuijken & Merrifield 1995; O’Neill & Dubinski 2003; Kormendy & Kennicutt 2004; Athanassoula 2005). Depending on, e.g., details of the equation of state, feedback, and rotational support of the gas disk, it may also amount to steady disk contraction (Debattista et al. 2006) or emergence of a two-component disk (Kaufmann et al. 2007; Foyle et al. 2008). A number of effects will be important in this regime, including the effects of feedback in pressurizing the disk and smoothing out substructure, and the role of accretion and mergers in rebuilding the disk as such evolution continues (since it is occurring on timescales \(\sim a\) several Gyr, comparable to the characteristic timescales for new accretion and mergers).

These effects make it difficult to predict the net effect of such evolution. For example, if in a pure gas merger of mass ratio \(\mu\) the specific angular momentum (on average) is increased (by addition of specific angular momenta plus orbital angular momentum) by an amount \(\sim \epsilon_m H_{\text{disk}}\), but then the induced bar (of amplitude \(\sim \mu\)) loses its angular momentum \(\sim \mu J_{\text{disk}}\) on a timescale \(N t_{\text{rot}}\), the sequence of mergers and induced bars compete (given the cosmologically expected timescale \(\approx \mu J_{\text{disk}}\) on a timescale \(N t_{\text{rot}}\), the sequence of mergers and induced bars compete {\(\epsilon_{\text{disk}} H_{\text{disk}} [\epsilon_m - \mu/N m_f]\)}—i.e., more major events will tend to lead to angular momentum loss in gas, whereas the net effect of very minor mergers and
smooth gas accretion, even where it induces instabilities, may be to “spin up” the disks).

As discussed in Section 4.3.3, there is also an interesting regime of parameter space, namely minor mergers with mass ratios \( \sim 1:20–1:10 \) or so, in which the characteristic merger timescales and secular/internal evolution timescales are comparable. The secondary may be large enough to induce a significant bar/nonaxisymmetric response, but the merger/dynamical friction time may be sufficiently long that the primary could respond almost as if in isolation for several orbital periods. In such a case it becomes less clear whether the merger or the secular response of the disk is ultimately the dominant driver of evolution (and the answer probably depends on, e.g., the exact orbital parameters and stability properties of the disk, and may be sensitive to the feedback, the gas phase structure and pressure support, and the detailed halo structure). In any event, it is clear that these processes require study in a more complete cosmological context, and can contribute significantly to the bulge population (especially in less bulge-dominated galaxies, below the typical thresholds we simulate) over a Hubble time of evolution. However, although the bar itself may be triggered in the merger, the nature of the relative strength of the interaction and characteristic timescale for angular momentum loss make it not a violent process associated with the merger itself, but rather a secular process that should be considered more analogous to bars in nonmerging systems.

5. APPLICATION TO SEMIANALYTIC MODELS

Our results clearly have potential uses as prescriptions for analytic and semianalytic models of galaxy formation. Here, we summarize and give some simple recommendations for these applications.

When a merger is identified in a semianalytic model, the two key quantities we can predict here are the mass fraction of the disk that is destroyed (violently relaxed into a bulge) and the fraction of the cold gas in the disk that will lose angular momentum and contribute to the bulge by forming a compact starburst.

First, the stellar disks: in a merger of secondary mass \( M_2 \) with primary mass \( M_1 \), the secondary is destroyed (adding \( M_2 \) to the bulge) and the mass within a radius enclosing \( \approx M_2 \) in the primary is violently relaxed. If the primary were pure disk, this would add \( 2M_2 \) to the bulge. However, one can imagine the limit where the primary is entirely bulge-dominated inside that radius (with the stellar disk dominant only at much larger radii)—then the violent relaxation of the merger will act primarily to heat existing bulge stars, and only a mass \( M_1 \) will be added to the bulge. Obviously, its also true that if the total disk mass of \( M_1 \) is less than \( M_2 \), then that is a maximum to how much can be added to the bulge (i.e., really \( \text{MIN}(M_2, f_{\text{disk}} M_1) \) is added). For most purposes, this factor 2 possible range is not critical in the semianalytic models, and picking a constant (effective mean) fraction \( (0–1) \times M_2 \) to violently relax in the primary in all mergers is acceptable. However, if more detail is desired, an estimate of the mass profiles of bulge plus disk components in the primary can be used to determine the total primary disk mass within a radius enclosing a mass \( \approx M_2 \), and then that will be the fraction violently relaxed. For a Hernquist (1990) bulge and exponential disk obeying roughly the observed size–mass relations from Shen et al. (2003), the primary disk mass that should be violently relaxed in a merger with mass \( M_2 \) can be approximated as \( f_{\text{disk}} \equiv M_2/(1 + [M_1/M_2]^\alpha) \), where

\[
\mu \equiv M_2/M_1,
\]

\( \theta \) and \( b \), equivalently the orbital parameters of the merger. As discussed in Section 4.3.2, for cases that will merge the appropriate limit is \( b \rightarrow 0 \), since most of the action will occur on the final merging passages after the angular momentum is removed. We discuss what should be adopted for the orbital inclination \( \theta \) below.

5. \( R_4 \), the scale length of the disk stars.

In any semianalytic or analytic model, variables (1)–(3) should be well known beforehand. Given some choice of orbital
parameters and an assumed mass distribution of the disk, it is trivial then to translate Equation (7) into a fraction of the gas that will burst. Because orbital parameters are generally undetermined in these models, there are two choices for the assumed orbital inclination \( \theta \). First, one could draw a random value of \( \theta \) for each merger (uniformly sampling in \( \cos(\theta) \) as appropriate for an isotropic orbit distribution), and use Equations (9) and (10) for each merger. Alternatively, we can average over a random distribution of orbits and quote an “effective” orbital dependence \( F(\theta, b) \) for Equation (9). Note that this is only strictly appropriate if all disks have the same mass profiles and those are such that the enclosed mass is linear in \( R/R_* \) (otherwise the appropriate average would have to be weighted by other terms such as \( (1 - f_{\text{gas}}) \) in Equation (7)). In any case doing so yields an effective mean orbital dependence \( F(\theta, b) \approx 1.2 \).

The only remaining issue is the assumed mass profile of the disk. Here, models have some freedom. As we have emphasized, the exact profile (e.g., choice of exponential disk or some other profile) does not have a dramatic effect. What is important, however, is the assumption of how the gas is distributed relative to the stars. Recall, Equation (7), with the variables above inserted, gives that the gas inside some radius \( R_{\text{gas}} = x R_d \) (where \( x \) is a constant depending on those variables, and \( R_d \) is a characteristic scale length of the stellar disk) should lose angular momentum and participate in the burst. Given a gas mass profile \( M_{\text{gas}}(R/R_*, \text{gas}) \), in terms of a characteristic gas disk scale length \( R_{\text{gas}} \), this gives the gas mass that bursts, \( M_{\text{gas}}(R/R_*, \text{gas}) \). For our simulations, we have generally assumed (and can see that it is a good approximation) that the gas and stellar disks initially trace one another (\( R_d \approx R_{\text{gas}} \)). However, since our derivation and Equation (7) show that it is the gas inside some fraction of the stellar disk half-mass radius \( R_d \) that loses angular momentum, then if the gas is, e.g., much more extended than the stars, a lower gas fraction will end up in the burst. We discuss this in Section 4.3.5, and consider how such situations may in fact arise owing to, e.g., supernova feedback blowing gas out to large radii. Semianalytic models therefore have some freedom in adopting these prescriptions based on their implicit assumptions about feedback and disk formation, encapsulated effectively in our prescriptions as the ratio of the stellar to gas disk scale lengths \( R_d/R_{\text{gas}} \). Lacking some detailed model for both values in the semianalytic models, a constant value \( \approx 1 \) is probably a good choice (with the exact choice reflecting implicit assumptions about feedback and outer disk formation).

Those prescriptions define both the violently relaxed and starburst components induced in mergers of arbitrary mass ratios, gas content, and orbital parameters. If desired, appropriate scatter (a factor of \( \approx 2 \)) can be added to both quantities, reflecting the scatter we see between various numerical realizations (although it should still be ensured that, with scatter, the implied violently relaxed and burst fractions are within the sensible physical limits).

Although not discussed here, in Hopkins et al. (2008a, 2008g, 2008h, 2008e, 2008d) we consider how the sizes and velocity dispersions of these components should scale, and we refer to those papers for detailed analysis of those results. Briefly, we note that in the absence of dissipation, it is straightforward to calculate the size of the dissipationless component (the violently relaxed stars from the premerger stellar disk), given phase space and energy conservation. Roughly, this implies that the component will have the same (modulo projection effects since it transforms from a disk to a sphere) scale radius as the disk (or radius within the disk) from which it forms. Again, conservation of energy in subsequent dissipationless remergers, along with the assumption of preserved profile shape (which we demonstrate is reasonable in Hopkins et al. 2008g) yields the evolution in subsequent events of these radii (in a remerger of masses \( M_1 \) and \( M_2 \), the dissipationless bulge component will have final size \( R_f/R_1 \approx (1 + \mu^2)/(1 + \mu^2 R_1/R_2) \)). Dissipation complicates this—it is possible to solve separately for the size of the dissipational component by allowing for energy loss in the collision followed by (after angular momentum loss) collapse to a self-gravitating limit, and then subsequently evolve the component as a dissipationless body, added with the violently relaxed components to give a total bulge effective radius. Fortunately, Covington et al. (2008) perform such an exercise and we show in Hopkins et al. (2008a) that their results can be conveniently approximated (in both an analytic manner and as a fit to the results of numerical simulations) by the scaling: \( R_f(\text{bulge}) = R_1(f_{\text{sh}} = 0)/(1 + f_{\text{sh}}/f_0) \), where \( f_{\text{sh}} \) is the total mass fraction of the bulge/spheroid which originally formed dissipationally (as opposed to being violently relaxed). \( R_1(f_{\text{sh}} = 0) \) is the radius the system would have if purely dissipationless (calculated as described above), and \( f_0 \approx 0.25–0.35 \) is a constant.

Our modeling could also be applied in the manner described in Section 4.3.4 to BH (nonmerging) encounters, but we caution that these are usually ill-defined in semianalytic models (and if adopted, the cautions in Section 4.3.4 about the appropriate meaning of the impact parameter adopted should be borne in mind). In any case, the rapid suppression of bursts with increasing impact parameter means that such cases should be relatively unimportant in a representative cosmological ensemble.

6. DISCUSSION AND CONCLUSIONS

We have derived a general physical model for how disks survive and/or are destroyed in mergers and interactions. Our model describes both the dissipational and dissipationless components of the merger, and allows us to predict, for a given arbitrary encounter, the stellar and gas content of the system that will be dissipationlessly violently relaxed, dissipationally lose angular momentum and form a compact central starburst, or survive (without significant angular momentum loss or violent relaxation) to re-form a disk. We show that, in an immediate (short-term) sense, the amount of stellar or gaseous disk that survives or re-forms following a given interaction can be understood purely in terms of simple, well understood gravitational physics. Knowing these physics, our model allows us to accurately predict the behavior in full hydrodynamic numerical simulations across as a function of the merger mass ratio, orbital parameters, premerger cold gas fraction, and mass distribution of the gas and stars, in simulations which span a wide range of parameter space in these properties as well as prescriptions for gas physics, stellar and AGN feedback, halo and initial disk structural properties, redshift, and absolute galaxy masses.

The fact that we can understand the complex, nonlinear behavior in mergers with this analytic model, and moreover that (for given conditions at the time of merger) our results are independent of the details of prescriptions for gas physics, star formation, and feedback, owes to the fact that the processes that strip angular momentum from gas disks and violently relax stellar disks are fundamentally dynamical.

Gas, in mergers, primarily loses angular momentum to internal gravitational torques (from the stars in the same disk)
Figure 16. Summary of our comparison between simulations and analytic model for the mass of disks in merger remnants as a function of appropriate orbital parameters, merger mass ratio, and premerger cold gas content. We plot our model prediction vs. the simulation remnant disk fraction for all \( \sim 400 \) full hydrodynamic merger simulations considered in this paper (shown in both a linear and logarithmic scale). Symbols encode some of the parameter studies we consider: orbital parameters, galaxy masses, initial merger redshift, choice of feedback prescription, merger mass ratio, and presence or absence of BHs, as labeled. For each subset of simulations, we sample a wide range in initial and premerger gas fractions \( f_{\text{gas}} = 0-1 \). The solid line is a one-to-one relation. In all cases, our predictions agree well with the simulations, with no systematic offsets owing to any of the parameters we have varied. At high \( f_{\text{disk}} \), our predictions are accurate to an absolute uncertainty of \( \sim 0.05-0.10 \) in \( f_{\text{disk}} \). At low \( f_{\text{disk}} \lesssim 0.1 \), our predictions are accurate to a factor of \( \sim 2-3 \) (down to \( f_{\text{disk}} \lesssim 1\% \), where it is difficult to reliably identify disks in the remnant).

(A color version of this figure is available in the online journal.)

owing to asymmetries in the galaxy induced by the merger (on the close passages and final coalescence of the secondary, during which phase the potential also rapidly changes, scattering and violently relaxing the central stellar populations of the stellar disk).\(^4\) Hydrodynamic torques and the direct torquing of the secondary are second-order effects, and inefficient for all but pathological orbits.

Once gas is efficiently drained of angular momentum, there is little alternative but for it to fall to the center of the galaxy and form stars, regardless of the details of the prescriptions for star formation and feedback—we show that even strong supernova-driven winds (with mass loading efficiencies several times the SFR and wind mass-loading velocities well above the halo escape velocity) do not significantly effect our conclusions. Such processes, after all, can blow out some of the gas, but they cannot fundamentally alter the fact that cold gas with no angular momentum will be largely unable to form any sort of disk, or the fact that a galaxy’s worth of gas compressed to high densities and small radii will inevitably form a large mass in stars.

For these reasons, many processes and details that are important cosmologically (systematically changing, e.g., the premerger disk gas fractions)—in some sense setting the initial conditions for our idealized study of what happens in mergers—do not alter the basic dynamical behavior within the mergers themselves, and therefore do not change our conclusions.

Figure 16 summarizes our results for the ensemble of our simulations. We compare the fraction of the baryonic galaxy mass in the merger remnant that is in a surviving postmerger disk to that predicted by our simple model scalings, and find good agreement over the entire range in disk and bulge mass fractions sampled, with surprisingly small scatter given the complexity of behavior in mergers. We highlight several of the parameter studies, showing that—for fixed mass ratio, orbital parameters, and gas content at the time of the final merger, none of these choices systematically affect our predictions (note that these are not the only parameters varied—the complete list is discussed in Section 2, but it is representative). That is not to say they cannot affect them indirectly, by, e.g., altering how much gas is available at the time of merger—but it emphasizes

\(^4\) We again note that although we have described these asymmetries as “bars” or “barlike” at certain points in this paper, there are a number of properties of the nonaxisymmetric distortions induced in mergers (discussed in Sections 4.3.2 and 4.3.7) that make them—at least over the short relaxation timescale of the merger—dynamically distinct from traditional bar instabilities in isolated systems.
that the processes we model and use to form our predictions, the processes that dominate violent relaxation and the loss of angular momentum in gas in mergers, are fundamentally dynamical.

This allows us to make robust, accurate physical predictions independent of the (considerable) uncertainty in feedback physics and subresolution physics of the ISM. Regardless of how those physics alter the “initial” conditions, they do not change basic dynamical processes, and so do not introduce significant uncertainties in our model.

In turn, this means that we can use our model to understand just why and how feedback is important for the cosmological survival of disks. Why, in short, have various works (see Springel & Hernquist 2005; Robertson et al. 2006a; Governato et al. 2007) concluded that strong feedback is essential for enabling disk survival in mergers? Our results show that it is not that feedback somehow makes the disk more robust to the dynamical torques within the merger, in any instantaneous sense. These torques, at least within the critical radii where the gravitational perturbation from the merger is large and in resonance, are sufficiently strong that any reasonable feedback prescription is a dynamically negligible restoring force. Rather, feedback has two important effects that fundamentally alter the conditions in the merger: first, it allows the galaxy to retain much higher gas content going into the merger. Without feedback from, e.g., star formation and supernovae contributing to heating and pressurizing the ISM and redistributing gas spatially, isolated gas-rich disks may be unstable to fragmentation. Even if fragmentation is avoided, it is well known that star formation in simulations proceeds efficiently under these conditions. This would leave the disks essentially pure stars (even for idealized simulations beginning with \( \sim 100\% \) gas disks; see Springel et al. 2005b) by the time of the merger, which guarantees that a major merger will inevitably violently relax the stars (this is a simple collisionless mixing process, and under such circumstances is inescapable). With large gas fractions, however, the system relies on stripping angular momentum from the gas to form new bulge stars, which in turn relies on internal torques from induced asymmetries in the stellar disk. If the gas fractions are sufficiently large, there is little stellar disk to do any such torquing, and the gas survives largely intact.

Second, feedback from supernovae and stellar winds moves the gas to large radii, where it does not feel significant torques from the merger. Again, recall that the most efficient torquing is driven by the internal stellar disk of the galaxy, and as such is most efficient at torquing gas within small radii (this can be thought of as analogous to the well known corotation condition for isolated disk bars). If star formation-driven feedback has blown much of the gas to large radii, then there is little gas inside the radius where torques can efficiently strip angular momentum, yielding little induced starburst and largely preserving the gas disk at large radii.

Not only can we qualitatively identify these requirements for feedback processes, but we can more precisely use our model to set quantitative limits on how much gas must be retained and/or the radii it must be redistributed to in order to enable disk survival under various conditions. This also clearly implies that disks must be able to avoid fragmentation and strong local gravitational instabilities when they achieve these gas fractions. This provides a valuable constraint for feedback models—how those models affect star formation efficiencies, the “blowout” of gas, and the local hydrodynamic state (effective equations of state and phase structure) of ISM gas—and should be useful for calibrating their (still largely phenomenological) implementations in both numerical and semianalytic models of galaxy formation.

Our predictions are also of interest in any cosmological model for the emergence of the Hubble sequence, since they apply not just to disk-dominated galaxies but to small disks in bulge-dominated systems. We give a number of simple prescriptions for application of our conclusions to analytic and semianalytic models of galaxy formation, which can be used to predict the distribution of bulge to disk ratios in cosmological ensembles. But even without reference to a full such model, a number of interesting consequences are immediately apparent.

First, it is a well known problem that theoretical models systematically overpredict the abundance and mass fractions of bulges in (especially) low-mass galaxies. This is true even in, e.g., semianalytic models, which are not bound by resolution requirements and can adopt a variety of prescriptions for behavior in mergers. However, it is also well established observationally that disk gas fractions tend to be very high in this regime, with large populations of gas-dominated disks at \( M_* < 10^{10} M_\odot \) (Bell & de Jong 2001; Kannappan 2004; McGaugh 2005). Our models predict that bulge formation should, therefore, be strongly suppressed in precisely the regime required by observations. For, e.g., disks with \( M_* < 10^{9} M_\odot \) where observations suggest typical gas fractions \( \sim 60\%–80\% \), our results show that even a 1:1 major merger would typically yield a remnant with only \( \sim 30\% \) bulge by mass—let alone a more typical 1:3–1:4 mass-ratio merger, which should yield a remnant with \( < 20\% \) bulge. That is not to say that it is impossible to form a bulge-dominated system at these masses, but it should be much more difficult than at high masses, requiring either unusually gas-poor systems, violent merger histories, or rarer merging orbits that are more efficient at destroying disks.

Our conclusions therefore have dramatic implications for the abundance of bulges and typical morphologies and bulge-to-disk ratios at low galaxy masses and in gas-rich systems. Low-mass systems, when a proper dynamical model of bulge formation in mergers is considered, should have lower bulge-to-disk ratios—by factors of several, at least—than have been assumed and modeled in previous theoretical models. Whether this alone is sufficient to resolve the discrepancies with the observations remains to be seen, but it is clearly of fundamental importance that future generations of models incorporate this scaling.

Second, the importance of this suppression owing to gas content in disks will be even more significant at high redshifts. Observations suggest (see Erb et al. 2006b) that by \( z \sim 2 \), even systems with masses near \( \sim L_* \) (\( M_* \sim 10^{10}–10^{11} M_\odot \)) may have gas fractions as high as \( f_{gas} \sim 0.6 \). In this regime, the same argument as above should apply, dramatically suppressing the ability of mergers to destroy disks. Moreover, since most of the mass density is near \( L_* \), this can change not just the behavior in a specific mass regime but significantly suppress the global mass density of spheroids, modifying the predicted redshift history of bulge formation. (Note that this will not change very much when \textit{stars} form, so it has little or no effect on, e.g., the ages of \( z \sim 0 \) spheroids).

This redshift evolution may also explain the solution to a fundamental problem in reconciling the observed disk populations with CDM cosmologies. Integrated far enough back in time, every galaxy is expected to have experienced a significant amount of major merging. In extreme cases, the mass of the system when it had its last such merger may be so small that it
would not be noticed today, but in general, it does not require going far back in redshift (to perhaps \( z \sim 2–4 \) before almost every \( z = 0 \) galaxy should have had such a merger). How, then, can the abundance of systems with relatively small (or even no) visible bulges be explained? Our conclusions here highlight at least part of the answer: as you go back in time, the gas fractions of systems are also higher, nearing unity. So even though, integrating sufficiently far in time, every system has experienced major mergers, it is also true that the systems were increasingly gas-rich, and therefore that the impact of those mergers was more and more suppressed. Only mergers at later times, below certain gas fraction thresholds, will typically destroy disks.

Third, to the extent that bulge formation is suppressed at increasing redshifts, the existence of an \( M_{\text{BH}}-M_{\text{bulge}} \) relation (e.g., Magorrian et al. 1998) implies that BH growth should also be suppressed. Indeed, bulge formation is suppressed specifically because gas cannot efficiently lose angular momentum in mergers if the systems are gas-dominated—if the gas cannot lose angular momentum efficiently, then it certainly cannot efficiently be accreted by the nuclear BH. Since this pertains to gas on the scales of galactic disks, it is probably not relevant for the formation of “seed” BHs at very high redshift, but it will in general inhibit the growth of BHs owing to early merging activity. At the same time, of course, higher gas fractions in general imply increasing fuel supplies for BH growth, so the effects are not entirely clear, and more detailed models are needed to see how this impacts the history of BH growth and quasar luminosity functions. Nevertheless, this may in part explain why, above \( z \sim 2 \) (where, for the argument above, these effects become important for the global mass density of spheroids), the global rate of BH growth (i.e., total quasar luminosity density) appears to decline much more rapidly with increasing redshift than the SFR density (compare, e.g., Hopkins & Beacom 2006; Hopkins et al. 2008f, 2007d).

Fourth, our models imply that a large fraction of bulges and disks survive mergers together, rather than being formed entirely separately. It is often assumed that classical bulges—being similar to small ellipticals in most of their properties—were formed initially in major mergers, as entirely bulge-dominated systems, and then accreted new gaseous and stellar disks at later times. Although nothing in our modeling would prevent this from happening, our analytic and simulation results generally lead to the expectation that a large (perhaps even dominant) fraction of the bulge population did not form in this manner, but rather formed in situ from minor mergers or less efficient major mergers (in, e.g., very gas-rich systems). Observations tracing the evolution of disk components, kinematics, and morphology in the last \( \sim 10 \) Gyr increasingly suggest that such coformation or disk regeneration scenario is common (see Hammer et al. 2005; Conselice et al. 2005; Flores et al. 2006; Puech et al. 2008, and references therein). In short, a system with a mass fraction \( \sim 0.1–0.2 \) in a bulge could be the remnant of an early, violent major merger (when the system was \( \sim 0.1 \) times its present mass) with a reaccreted disk, or could be the remnant of a typical (low to intermediate gas fraction) \( 1:10–1:5 \) mass ratio minor merger, or could even be the remnant of a gas-rich major merger (mass ratio \( \lesssim 1:3 \), if \( f_{\text{gas}} \) is sufficiently large).

Based on a simple comparison of typical merger histories, we would actually expect that the minor merger mechanism should be most common, but all may be non-negligible. Fundamentally, the physics forming the bulge (torquing the gas within some radius owing to internal asymmetries and violently relaxing stars within a corresponding radius) are the same in all the three cases, and moreover other indicators such as their stellar populations will be quite similar (in all the cases, the bulge will appear old: this is both because the central stars in even present-day disks are much older than those at more typical radii, and because in any case star formation will cease within the bulge itself, as opposed to the ongoing star formation in the disk, and stellar population age estimates are primarily sensitive to the amount of recent or ongoing star formation; see Trager et al. 2000). This is also not to say that mergers are the only means of producing bulges. Secular evolution of, e.g., barred disks probably represents an increasingly important channel for bulge evolution in later-type and more gas-rich systems (see Christodoulou et al. 1995a; Sheth et al. 2003; Mayer & Wadsley 2004; Debattista et al. 2004; Jogee et al. 2004; Kormendy & Kennicutt 2004; Marinova & Jogee 2007), and may even be related (albeit through longer timescales of “isolated,” postmerger evolution and different physics) to initial bar formation or “triggering” in mergers. More detailed theoretical work and analysis of cosmological simulations is needed to develop observational probes that can distinguish between these histories.

Further work is specifically needed to investigate the processes at work in minor mergers with mass ratios \( \sim 1:10–1:5 \) (\( \mu \sim 0.05–0.1 \)), which cosmological simulations suggest are an important contributor to the growth of disks, especially in later-type systems (Maller et al. 2006; Fakhouri & Ma 2008; Stewart et al. 2008). In more minor mergers \( \mu \ll 0.1 \), the secondaries are sufficiently small and dynamical friction times sufficiently long that the disk is unlikely to feel significant external perturbations. More major mergers \( \mu \gtrsim 0.1 \), the cases of interest here, induce sufficiently large responses in the disk and evolve sufficiently rapidly that they can be considered “merger dominated” for the reasons in Sections 4.3.2 and 4.3.7. But in the intermediate regime, internal amplification of instabilities in a traditional secular fashion may occur on a timescale comparable to or shorter than the evolution of the secondary orbit, potentially leading to a more complex interplay between the two. It is not entirely clear whether such a system would remain “locked” to the driven perturbation, or function as a purely secular system (merely initially driven by the presence of the secondary), or some nonlinear combination of both. A more detailed comparison of the relevant timescales for these processes and their relation to, e.g., cosmological triggering of bars and large-scale nonaxisymmetric modes in disks will be the subject of future study (in preparation).

Our results are also of direct interest to models of spheroid formation in ellipticals and S0 galaxies. As discussed in Section 1, it is increasingly clear that embedded subcomponents—constituting surviving gaseous and stellar disks—are both ubiquitously observed and critical for theoretical models to match the detailed kinematics and isophotal shapes of observed systems (Naab et al. 2006; Cox et al. 2006a, 2006b; Robertson et al. 2006b; Jesseit et al. 2007; Hopkins et al. 2008a, 2008b). We have developed a model that allows us to make specific predictions for how disks survive mergers, including both the survival of some amount of the premerger stellar disks and the postmerger re-formation of disks and rotationally supported components from gas that survives the merger without losing most of its angular momentum.

Figure 16 shows that we can extend these predictions with reasonable accuracy to surviving rotational systems containing as little as \( \sim 1\% \) of the remnant stellar mass, comparable to small central subcomponents and subtle features giving rise to, e.g., slightly disky isophotal shapes (see Ferrarese et al.
1994; Lauer et al. 2005; McDermid et al. 2006). Owing to the combination of resolution requirements and desire to understand the fundamental physics involved, most theoretical studies of these detailed properties of ellipticals have been limited to idealized studies of individual mergers. Our results allow these to be placed in a more global context of cosmological models and merger histories. Moreover, our models allow the existence of such features (or lack thereof) to be translated into robust constraints on the possible merger histories and gas-richness of spheroid-forming mergers. Further, Hopkins et al. (2008a, 2008b), studied how the dissipational starburst components arising in gas-rich mergers are critical to explaining the observed properties and scaling relations of ellipticals, and how these components can both be extracted from and related to observed elliptical surface brightness profiles. Because both the starburst and surviving disks arise from gas in mergers, the combination of constraints from the central stellar populations, studied therein, with constraints on the survival and/or loss of gas angular momentum in mergers studied here, should be able to break some of the degeneracies in, e.g., premerger gas fractions and merger histories in order to enable new constraints and understanding of spheroid merger histories, and new tests of models for spheroid formation in gas-rich mergers.

These points relate to a number of potentially testable predictions of our models. These include the in situ formation of bulges from various types of mergers, and possible associated stellar population signatures, the presence of embedded disks in ellipticals, and how their sizes and mass fractions scale with, e.g., the masses and formation times of ellipticals (and how this relates to gas fractions and stellar populations in observed disks). In general, for similar merger histories, the increasing prevalence of later-type galaxies (S0’s and Sa’s) at lower masses where disks are characteristically more gas rich is a natural consequence of our predictions here, and it is straightforward to convert our predicted scalings into detailed predictions for the abundance and mass fractions of disks given some simplified merger histories. To the extent that these processes also give rise to disk heating and/or increasing velocity dispersions in disks, or changing kinematics in both disks and bulges, then there should be corresponding relationships between galaxy shapes, kinematics, and bulge-to-disk ratios along the Hubble sequence. We investigate these possible correlations and tests in subsequent papers (in preparation).

Altogether, our results here elucidate the relevant physics important for both dissipational and dissipationless bulge formation in mergers. They support a new paradigm in which to view bulge and disk formation: gas-richness is not simply a “tweak” to existing models of bulge formation and disk destruction in mergers. Rather, if disks are sufficiently gas rich, the qualitative character of mergers is different, with inefficient angular momentum loss giving rise to disk-dominated remnants. This process is not inherently governed by poorly understood feedback physics (although such feedback may be critical for establishing the conditions necessary in the first place), but rather by well understood gravitational physics, and as such is robust and fundamentally inescapable. Aspects of galaxy populations such as the continuum of relative bulge and disk mass ratios are not simply consequences of, e.g., different amounts of accretion, but can arise owing to the continuum in efficiencies of disk destruction as a function of merger mass ratios, orbital parameters, and gas content. The relative (lack of) abundance of bulges at low galaxy masses and high redshift is a basic consequence of the dynamics of how gas loses angular momentum in mergers, even for similar merger histories. In short, the baryonic physics of mergers ensures that, despite the near self-similarity of the physics and merger histories of their host halos, disk and bulge formation are not a self-similar process, influenced dramatically (well out of proportion to the absolute cold gas mass fractions) by the gas-richness of the baryonic systems.

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