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CONSEQUENCES OF GRAVITATIONAL RADIATION RECOIL

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ABSTRACT

Coalescing binary black holes experience an impulsive kick due to anisotropic emission of gravitational waves. We discuss the dynamical consequences of the recoil accompanying massive black hole mergers. Recoil velocities are sufficient to eject most coalescing black holes from dwarf galaxies and globular clusters, which may explain the apparent absence of massive black holes in these systems. Ejection from giant elliptical galaxies would be rare, but coalescing black holes are displaced from the center and fall back on a time scale of order the half-mass crossing time. Displacement of the black holes transfers energy to the stars in the nucleus and can convert a steep density cusp into a core. Radiation recoil calls into question models that grow supermassive black holes from hierarchical mergers of stellar-mass precursors.

Subject headings: black hole physics — gravitation — gravitational waves — galaxies: nuclei

1. KICK AMPLITUDE

In a companion paper (Favata, Hughes & Holz 2004; hereafter Paper I), the amplitude of the recoil velocity experienced by a binary black hole (BH) due to anisotropic emission of gravitational radiation during coalescence is computed. Here we explore some of the consequences of the kicks (Redmount & Rees 1989): the probability that BHs are ejected from galaxies and the implications for BH growth; the time scale for a kicked BH to return to the center of a galaxy; the effect of displacement on nuclear structure; and other observational signatures of the kicks. Unless otherwise indicated, notation is the same as in Paper I.

For inspiral from a circular orbit, the kick velocity is a function of the binary mass ratio $q = m_1/m_2 \leq 1$, the BH spins \tilde{a}_1 and \tilde{a}_2 , and the initial angle ι between the spin of the larger BH and the orbital angular momentum of the binary. Following Paper I, the spin of the smaller BH is ignored. Although Paper I only considers the cases $\iota = 0$ and $\iota = 180$, the recoil for arbitrary inclination is likely to be bounded between these extreme values. Also, the detailed inclination dependence is unimportant in comparison with the large uncertainty already present in the contribution to the recoil from the final plunge and coalescence. We will therefore assume that the restriction to equatorial-prograde/retrograde orbits ($\tilde{a}_2 = [-1, 1]$) considered in Paper I encompasses the characteristic range of recoil velocities.

Figure 2b of Paper I shows upper- and lower-limit estimates of the recoil velocity as a function of the effective spin parameter \tilde{a} for reduced mass ratio $\eta = \mu/M = q/(1+q)^2 = 0.1$. The *upper limit* for $\eta = 0.1$ is well fit in the range $-0.9 \leq \tilde{a} \leq 0.8$ by the following fifth-order

polynomial:

$$V_{\text{upper}} = 465 \text{ km s}^{-1} \frac{f(q)}{f_{\text{max}}} (1 - 0.281\tilde{a} - 0.0361\tilde{a}^2 - 0.346\tilde{a}^3 - 0.374\tilde{a}^4 - 0.184\tilde{a}^5). \quad (1)$$

Fitchett's (1983) scaling function $f(q)/f_{\text{max}}$, with $f(q) = q^2(1-q)/(1+q)^5$, equals 0.433 for $\eta = 0.1$. The *lower limit* curve of Paper I is well fit by

$$V_{\text{lower}} = 54.4 \text{ km s}^{-1} \frac{f(q)}{f_{\text{max}}} (1 + 1.22\tilde{a} + 1.04\tilde{a}^2 + 0.977\tilde{a}^3 - 0.201\tilde{a}^4 - 0.434\tilde{a}^5). \quad (2)$$

We convert these expressions into estimates of the bounds on V_{kick} as follows. First, as discussed in Paper I, there is an ambiguity in how one translates the physical spin parameter \tilde{a}_2 of the larger hole into the effective spin parameter \tilde{a} of equations (1) and (2). Here we adopt the Damour (2001) relation $\tilde{a} = (1 + 3q/4)(1+q)^{-2}\tilde{a}_2$. Second, Fitchett's scaling function assumes that both bodies are non-spinning, and vanishes when $q = 1$. In fact, when $\tilde{a} \neq 0$, significant recoil would occur even for $q = 1$ due to spin-orbit coupling. We can guess the approximate form of a new scaling function by examining the spin-orbit corrections (Kidder 1995) to Fitchett's recoil formula. For equatorial orbits, equation (4) of Paper I suggests that $f(q)$ should be multiplied by the factor $|1 + (7/29)\tilde{a}_2/(1-q)|/|1 + (7/29)\tilde{a}_2/(1-q')|$, where $q' = 0.127$ is the value used in defining V_{upper} and V_{lower} in equations (1) and (2).

Figure 1 plots upper and lower limits to V_{kick} as functions of \tilde{a}_2 and q . The average over \tilde{a}_2 of the upper limit estimates are $\sim (138, 444, 154) \text{ km s}^{-1}$ for $q = (0.1, 0.4, 0.8)$; Figure 1 suggests a weak dependence on \tilde{a}_2 . Lower limit estimates are more strongly spin-dependent; the averages over \tilde{a}_2 are $\sim (21.1, 63.6, 24.9) \text{ km s}^{-1}$ for the same values of q . For moderately large spins ($\tilde{a}_2 \gtrsim 0.8$) and prograde capture, the lower limit estimates exceed 100 km s^{-1} for $0.2 \lesssim q \lesssim 0.6$. In what follows, we will assume that $\sim 500 \text{ km s}^{-1}$ is an absolute upper limit to V_{kick} .

2. ESCAPE

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When $V_{\text{kick}} \geq V_{\text{esc}} \equiv \sqrt{2\phi(\mathbf{r}=0)}$, with $\phi(\mathbf{r})$ the gravitational potential of the system (galaxy, dark matter halo) hosting the BH, the BH has enough kinetic energy to escape. Figure 2 shows central escape velocities in four types of stellar system that could contain merging BHs: giant elliptical galaxies (E), dwarf ellipticals (dE), dwarf spheroidals (dSph) and globular clusters (GC). We fit the trend $\log(V_{\text{esc}}/1 \text{ km s}^{-1}) = \lambda - \beta M_V$ separately for each class of object. dEs and GCs each separately establish a relation $L \sim V_{\text{esc}}^2$; for GCs, this is compatible with the relation found by Djorgovski et al. (1997). The E sample is consistent with the Faber-Jackson (1976) relation.

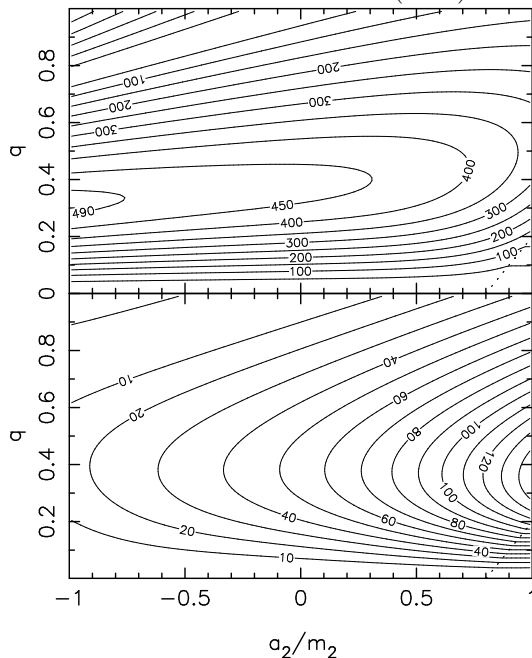


FIG. 1.— Upper limit (top) and lower limit (bottom) estimates of V_{kick} as functions of mass ratio q and spin of the larger black hole \tilde{a}_2 . Units are km s^{-1} . Values of \tilde{a}_2 and q corresponding to $\tilde{a} > 0.8$ lie in the region to the right of the dotted line. Since equations (1) and (2) are not valid for $\tilde{a} > 0.8$, \tilde{a} was replaced by 0.8 in this region.

The solid line in Figure 2 shows escape velocities from the dark matter (DM) halos associated with the luminous stellar systems. To relate halo properties to galaxy luminosities, we use the concordance Λ CDM model M1 of Yang, Mo & van den Bosch (2003). This function gives the number of galaxies in the luminosity range $L \pm dL/2$ that reside inside a halo with virial mass M_{vir} . The average luminosity L_1 of the brightest (“central”) galaxy in the halo of mass M_{vir} is implicitly given by the condition $\int_{L_1}^{\infty} \Phi(L|M_{\text{vir}})dL = 1$. Inverting this we obtain $M_{\text{vir}}(L_1)$ and relate this mass to the escape velocity via $V_{\text{esc}}^2 = 2cg(c)GM_{\text{vir}}/R_{\text{vir}}$ where R_{vir} is the virial radius of the halo, c is the concentration of a halo obeying the Navarro, Frenk & White (1996; hereafter NFW) profile, and $g(c) = [\ln(1+c) - c/(1+c)]^{-1}$ (e.g., Lokas & Mamon 2001). Both R_{vir} and c are functions of M_{vir} and the redshift z (e.g., Bryan & Norman 1998; Bullock et al. 2001). At $z = 0$ the average escape velocity is given by $V_{\text{esc}} = 239 \text{ km s}^{-1} (m_{11}/h)^{1/2}$, where $M_{\text{vir}} = 10^{11} m_{11} M_{\odot}$ and h is the Hubble parameter, set to 0.7 in Figure 2. The upturn in escape velocity for galaxies brighter than $M_V \sim -20$ is a consequence

of the increase in the occupation number of their host halos. The dashed line in Figure 2 shows the escape velocity from the combined luminous + DM potential for the E galaxies, using the scaling relation derived above to describe the luminous component.

Figure 2 suggests that the consequences of the kicks are strikingly different for the different classes of stellar system that might host BHs. Escape velocities from E galaxies are dominated by the stellar contribution to the potential; in the sample of Faber et al. (1997), $V_{\text{esc}} \gtrsim 450 \text{ km s}^{-1}$ even without accounting for DM. This exceeds even the upper limits in Figure 1. Hence, the kicks should almost never unbind BHs from E galaxies. The tight correlations observed between BH mass and bulge luminosity (McClure & Dunlop 2002; Erwin, Graham & Caon 2003) and velocity dispersion (Ferrarese & Merritt 2000; Gebhardt et al. 2000) could probably not be maintained if escape occurred with any significant frequency from luminous galaxies.

The existence of DM significantly affects the escape probability from dE and dSph galaxies, implying kicks of $\sim 300 \text{ km s}^{-1}$ and $\sim 100 \text{ km s}^{-1}$ respectively for escape. In the absence of DM, these numbers would be $\sim 100 \text{ km s}^{-1}$ and $\sim 20 \text{ km s}^{-1}$ respectively. Hence, kicks of order 200 km s^{-1} would unbind BHs from dSph galaxies whether or not they contain DM, while dE galaxies could retain their BHs if they are surrounded by DM halos.

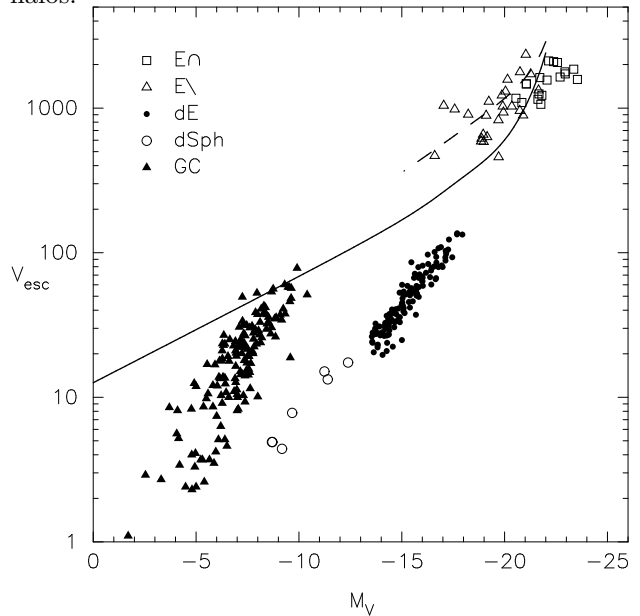


FIG. 2.— Central escape velocities in km s^{-1} in four types of stellar system that could harbor merging BHs. E galaxy data are from Faber et al. (1997), with separate symbols for core (\square) and power-law (\triangle) galaxies. dE data are from Binggeli & Jerjen (1998), with mass-to-light ratios from Mateo (1998). GC and dSph data are from the tabulation of Webbink (1996). Solid line is the mean escape velocity from the DM halos associated with the luminous matter. Dashed line is the escape velocity from the combined luminous + mean DM potentials for E galaxies.

Evidence for intermediate-mass black holes at the centers of galaxies fainter than $M_V \approx -19$ is sketchy (e.g., van der Marel 2003), although there is indirect (non-dynamical) evidence for BHs in faint Seyfert bulges (Filippenko & Ho 2003). We note that the dense nuclei

associated with BHs in galaxies like M32 ($M_V \approx -19$) become progressively less frequent at magnitudes fainter than $M_V \approx -16$ and disappear entirely below $M_V \approx -12$ (van den Bergh 1986). If the dense nuclei are associated with nuclear BHs (e.g., Peebles 1972), their absence could signal loss of the BHs via ejection. It is intriguing that these nuclei are sometimes observed to be displaced far from the galaxy center (Binggeli, Barazza & Jerjen 2000). Figures 1 and 2 imply that even kicks at the lower limits of Paper I would almost always unbind BHs from GCs.

3. EJECTION IN HIERARCHICAL MERGING SCENARIOS

The kicks have serious implications for models in which massive BHs grow from mergers of less massive seeds. In some of these models, the precursors are stellar- or intermediate-mass black holes produced in the collapse of the first stars (Population III) and the merging commenced in minihalos at redshifts as large as ~ 20 (Madau & Rees 2001; Volonteri, Haardt, & Madau 2003; Islam, Taylor & Silk 2003). We evaluate the plausibility of such models in light of the estimates of V_{kick} derived in Paper I. Kicks from gravitational wave emission may compete with high-velocity recoils (Saslaw, Valtonen, & Aarseth 1974) from (Newtonian) three-body interactions. While the Newtonian recoil occurs only when three BHs are present, which is contingent on the galaxy merger rate and the BH binary orbital decay rate, radiation recoil is present whenever BHs coalesce.

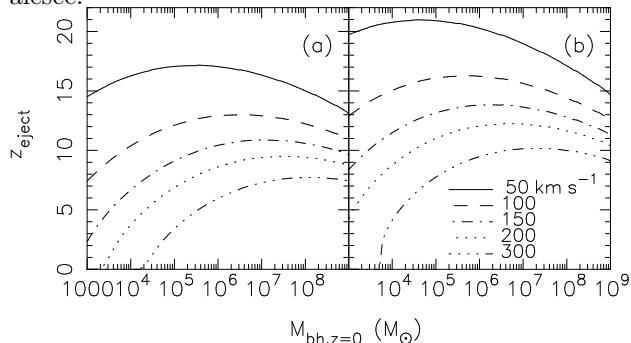


FIG. 3.— The maximum redshift z_{eject} at which (a) DM halos only, and (b) DM halos and the central galaxies combined, can confine BHs as a function of the $z = 0$ BH mass, for five values of the kick velocity. The depth of the galactic contribution to the potential was calculated by identifying the velocity dispersion of the stellar spheroid with the circular velocity of the halo (Ferrarese 2002).

The confining effect of DM halos in a hierarchical universe was smaller at higher redshifts when the average halo mass was smaller. We estimate the maximum redshift at which DM halos can confine the progenitors of the present-day BHs. Ferrarese (2002) derived a relation of the present-day BH mass $M_{\text{bh}}(z = 0)$ to the mass of the host halo $M_{\text{vir}}(z = 0)$. We use the Wyithe & Loeb (2002) form of the relation (their equation 11) to obtain the host halo mass and extrapolate the mass back in redshift via the accretion history model (Bullock et al. 2001) calibrated by Wechsler et al. (2002) on a set of numerical simulations of DM clustering in a Λ CDM universe. The accretion trajectory $M_{\text{vir}}(z) \propto e^{-\alpha z}$, where α is itself a function of the halo mass at $z = 0$, can be interpreted as the mass of the most massive, and thus the most easily confining parent halo at redshift z . We can then calculate

the escape velocity $V_{\text{esc}}(z)$ of the most massive progenitor halo as a function of redshift. Finally, we solve for z_{eject} such that $V_{\text{kick}} = V_{\text{esc}}(z_{\text{eject}})$; this is the maximum redshift at which the progenitors of the present-day BHs could have started merging. We also modelled the effect on z_{eject} of including the potential due to a stellar component, idealized as an isothermal sphere with core radius $r_h = 2GM_{\text{bh}}/\sigma^2$ and outer cutoff $10^3 r_h$; the 1D stellar velocity dispersion σ was related to halo circular velocity as in Ferrarese (2002).

The results for five representative choices of V_{kick} are shown in Figure 3. For $V_{\text{kick}} \sim 150 \text{ km s}^{-1}$, we find $z_{\text{eject}} < 11(14)$ over the entire range of M_{bh} ; the latter value is from the models that include a stellar component. For $V_{\text{kick}} \sim 300 \text{ km s}^{-1}$, the assembly of a $10^8 M_{\odot}$ BH must have started at $z \lesssim 8(10)$. Models that grow supermassive BHs from mergers of seeds of much lower mass at redshifts $z \gtrsim 10$ are thus disfavored due to the difficulty of retaining the kicked BHs. The effects of the kicks could be mitigated if early growth of the BHs was dominated by accretion, or if BH mergers were delayed until their halos had grown much more massive.

4. FALLBACK TIMES

A BH that has been kicked from the center of a stellar system with a velocity less than V_{esc} falls back and its orbit decays via dynamical friction against the stars and gas. We define the fallback time T_{infall} as the time required for a BH to return to a zero-velocity state after being ejected. The velocity with which the BH is ejected from the site of the merger is $V_{\text{eject}} = (M_{\text{bh}}/M_{\text{eff}})V_{\text{kick}} < V_{\text{kick}}$; here $M_{\text{eff}} = M_{\text{bh}} + M_{\text{bound}}$ with M_{bound} the mass in stars that remain bound to the BH after it is kicked. For recoil in a singular isothermal sphere nucleus $\rho \propto r^{-2}$, $M_{\text{eff}}/M_{\text{bh}} \approx (1.9, 1.5, 1.05, 1.00)$ when $V_{\text{kick}}/\sigma = (0.5, 1, 2, 3)$ where σ is the 1D stellar velocity dispersion; $M_{\text{bound}}/M_{\text{bh}} \propto (V_{\text{kick}}/\sigma)^{-4}$ for $V_{\text{kick}} \gg \sigma$.

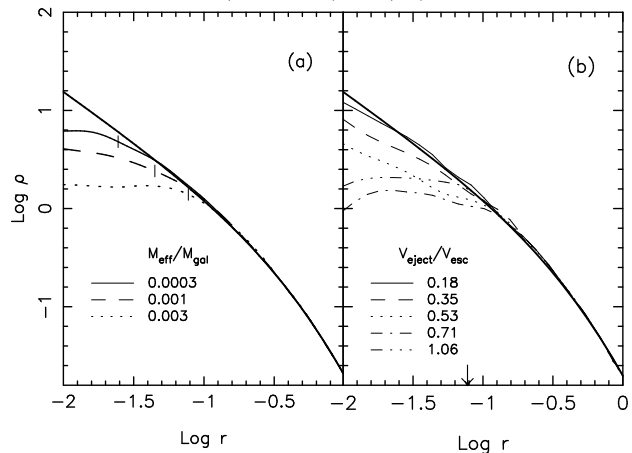


FIG. 4.— Effect on the nuclear density profile of black hole ejection. The initial galaxy model (black line) has a $\rho \propto r^{-1}$ density cusp. (a) Impulsive removal of the black hole. Tick marks show the radius of the black hole's sphere of influence r_h before ejection. A core forms with radius $\sim 2r_h$. (b) Ejection at velocities less than V_{esc} . The black hole has mass $0.003 M_{\text{gal}}$; the galaxy is initially spherical and the black hole's orbit remains nearly radial as it decays via dynamical friction. The arrow in panel marks r_h .

We evaluated T_{infall} for BHs kicked from the centers of Dehnen (1993)-law galaxies for which the central density obeys $\rho \propto r^{-\gamma}$. Bright E galaxies have $0 \lesssim \gamma \lesssim 1$

(Gebhardt et al. 1996), and cusps steeper than this are likely to be softened by the binary BH prior to coalescence (Milosavljević & Merritt 2001) and by the ejection itself (§5). Given values for M_{eff} and V_{eject} , the fall-back time in a spherical galaxy is given by the orbit-averaged dynamical friction equation (Cohn & Kulsrud 1978). For $V_{\text{eject}}/V_{\text{esc}} \lesssim 0.6$, infall times were found to be well approximated by $T_{\text{infall}} \approx T_{1/2}(V_{\text{eject}}/V_{\text{esc}})^{2.5(1+\gamma)}$ for $M_{\text{eff}} = 0.001M_{\text{gal}}$, where the period $T_{1/2}$ of a circular orbit at the galaxy’s half-light radius is given in terms of the galaxy’s visual luminosity by $T_{1/2} \approx 2 \times 10^8 \text{ yr} (L_V/10^{11} L_{\odot})^{1/2}$ (Valluri & Merritt 1998). Thus, return of a BH to a stationary state requires of order a few times 10^8 yr or less over a wide range of cusp slopes and galaxy luminosities for $V_{\text{eject}} \lesssim V_{\text{esc}}/2$. As indicated in Figure 2, this is the likely situation in the bright E galaxies. Infall times are especially short for $\gamma \geq 1$, since the BH experiences a strong impulsive frictional force as it passes repeatedly through the dense center. When $V_{\text{eject}} \lesssim \sigma$, the BH never moves far from its central position and it carries much of the nucleus with it. We carried out N -body simulations of this regime and found that return to zero velocity occurs in roughly one orbital period when $V_{\text{eject}} \lesssim \sigma$. In fainter dE and dSph galaxies, ejection would more often occur near V_{esc} and infall times could be arbitrarily long, determined primarily by the mass distribution at large radii.

In a nonspherical galaxy, an ejected BH does not pass precisely through the dense center on each return, delaying the infall. To test the effect of non-spherical geometries on the infall time, we carried out experiments in the triaxial generalizations of the Dehnen models (Merritt & Fridman 1996). Results were found to depend only weakly on the axis ratios of the models. Decay times in the triaxial geometry exhibit a spread in values depending on the initial launch angle, bounded from below by the decay time along the short axis. We found a mean at every $V_{\text{eject}}/V_{\text{esc}}$ that is $\sim 3 - 5$ times greater than in a spherical galaxy with the same cusp slope.

5. OBSERVABLE CONSEQUENCES OF THE DISPLACEMENT

Displacement of the BH also transfers energy to the nucleus and lowers its density within a region of size $\sim r_h$, the radius of the BH’s sphere of influence (defined here as the radius of a sphere containing a mass in stars equal to twice that of the BH). The simplest case to consider is $V_{\text{eject}} \gtrsim V_{\text{esc}}$; the BH and its entrained mass depart the nucleus on a time scale that is of order the

crossing time at r_h or less and do not return. The effect on the nucleus can be approximated by constructing a steady-state model of a galaxy containing a central point mass, then removing the point mass instantaneously and allowing the remaining particles to relax to a new steady state. Figure 4a shows the results for three values of $M_{\text{eff}}/M_{\text{gal}}$. Initial conditions consisted of 10^6 particles representing stars in a $\gamma = 1$ Dehnen model. We find that a core of roughly constant density forms within a radius of $\sim 2r_h$. Setting $\gamma = 2$ (not shown) results in a core of size $\sim r_h$. Figure 4b shows the change in the nuclear density profile for simulations with $V_{\text{eject}} < V_{\text{esc}}$. Significant changes in the central density require $V_{\text{eject}} \gtrsim 0.25V_{\text{esc}}$. We conclude that the recoil could affect the observable structure of nuclei, since radii of $\sim 2r_h$ are resolved in many nearby galaxies (Merritt & Ferrarese 2001).

The “mass deficits” seen at the centers of bright galaxies (Milosavljević et al. 2002; Ravindranath, Ho & Filippenko 2002) may be due to the combined effects of slingshot ejection and BH displacement, although we note that the large cores observed in some bright galaxies could probably not be produced by either mechanism (Milosavljević et al. 2002).

The X-shaped radio sources associated with giant E galaxies (Leahy & Parma 1992) are plausible sites of recent BH coalescence (Merritt & Ekers 2002). Displacement of the merged BHs from the galaxy center prior to ignition of the “active” lobes would imply a distortion of the X-morphology, in the sense that the “wings” (the inactive lobes) would be non-collinear near the center of the X. Such distortions are in fact a common feature of the X-sources (Gopal-Krishna, Biermann & Wiita 2003), although the linear scale of the distortions in some of the X-sources (e.g. $\sim 10 \text{ kpc}$ in NGC 236; Murgia et al. 2001) suggests that orbital motion of the merging galaxies may be a more likely explanation (Balcells et al. 1995).

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