MASSIVE PERTURBER–DRIVEN INTERACTIONS BETWEEN STARS AND A MASSIVE BLACK HOLE

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1. INTRODUCTION

There is compelling evidence that massive black holes (MBHs) lie in the centers of all galaxies (Ferrarese & Merritt 2000; Gebhardt et al. 2003; Shields et al. 2003), including the center of our Galaxy (Eisenhauer et al. 2005; Ghez et al. 2005). The MBH affects the dynamics and evolution of the galaxy’s center as a whole (e.g., Bahcall & Wolf 1976), and it also strongly affects individual stars or binaries that approach it. Such close encounters, which may be extremely energetic or involve non-gravitational interactions or post-Newtonian effects, have been the focus of many studies (see review by Alexander 2005). These processes include the destruction of stars by the MBH, either by falling whole through the event horizon, or by being first tidally disrupted and then accreted (e.g., Rees 1988); tidal scattering of stars on the MBH (Alexander & Livio 2001); the capture and gradual inspiral of stars into the MBH, accompanied by the emission of gravitational waves or by tidal heating (e.g., Alexander & Hopman 2003; Alexander & Morris 2003); or dynamical exchange interactions in which incoming stars or binaries energetically eject a star tightly bound to the MBH and are captured in its place very near the MBH (e.g., Alexander & Livio 2004; Gould & Quillen 2003).

The interest in such processes is driven by their possible implications for the growth of MBHs, for the orbital decay of a MBH binary, for the detection of MBHs, for gravitational wave (GW) astronomy, as well as by observations of unusual stellar phenomena in our Galaxy; for example, the puzzling young pop-

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2005; Berczik et al. 2006; or MBH feeding: Zhao et al. 2002; Miralda-Escudé & Kollmeier 2005) or in the hope that they may lead to significantly higher event rates for close encounter processes. These mechanisms include chaotic orbits in triaxial potentials (Norman & Silk 1983; Gerhard & Binney 1985; Merritt & Poon 2004; Holley-Bockelmann & Sigurdsson 2006) (the presence of a MBH may, however, destroy the triaxiality near the center; Merritt & Quinlan 1998; Holley-Bockelmann et al. 2002; Sellwood 2002); increased fraction of low angular momentum orbits in nonspherical potentials (Magorrian & Tremaine 1999; Berczik et al. 2006); accelerated resonant relaxation of angular momentum near the MBH where the orbits are Keplerian (Rauch & Tremaine 1996; Rauch & Ingalls 1998; Hopman & Alexander 2006a; Levin 2007); and perturbations by a massive accretion disk or an intermediate-mass black hole (IMBH) companion (Polnarev & Rees 1994; Zhao et al. 2002; Levin et al. 2005). Most of these mechanisms require special circumstances to work (e.g., specific asymmetries in the potential) or are short-lived (e.g., the IMBH will eventually coalesce with the MBH).

Here we explore another possibility that is more likely to apply generally: accelerated relaxation and enhanced rates of close encounters driven by massive perturbers (MPs). Efficient relaxation by MPs was first suggested by Spitzer & Schwarzschild (1943) as the time for a change of order unity in the phase space due to scattering, both in the context of scattering of Oort cloud comets to the Sun (Hills 1981; Bailey 1983) and the scattering of stars to a MBH in a galactic nucleus (Zhao et al. 2002). Zhao et al. (2002) suggested MPs as a mechanism for establishing the $M_{\text{MBH}} - \sigma$ relation (Ferrarese & Merritt 2000; Gebhardt et al. 2000) by fast accretion of stars and dark matter. They also noted the possibility of increased tidal disruption flares and accelerated MBH binary coalescence due to MPs. In this study we investigate in detail the dynamical implications of loss cone refilling by MPs. We evaluate its effects on the different modes of close interactions with the MBH, in particular three-body exchanges, which were not considered by Zhao et al. (2002), and apply our results to the Galactic center (GC), where observations indicate that dynamical relaxation is very likely dominated by MPs.

This paper is organized as follows. In § 2 we present the main concepts and procedures of our calculations. The observational data and theoretical predictions about MPs in the inner ∼100 pc of the GC are reviewed in § 3. In § 4 we explore the implications of relaxation by MPs for various types of interactions with the MBH. We summarize our results in § 5.

2. LOSS CONE REFILLING

In addition to stars, galaxies contain persistent dense structures such as molecular clouds, open clusters, and globular clusters with masses up to $10^4 - 10^7 M_\odot$. Such structures can perturb stellar orbits around the MBH much faster than two-body stellar relaxation (hereafter “stellar relaxation”), provided that they are numerous enough. This condition can be quantified by considering a test star randomly scattered by perturbers with masses in the interval $(M_p, M_p + dM_p)$ and number density $(dN_p/dM_p)dM_p$, approaching it with relative velocity $v$ on orbits with impact parameters in the interval $(b, b + db)$. The minimal impact parameter still consistent with a small-angle deflection is $b_{\text{min}} = 2GM_p/v^2$ (the capture radius), where $v$ is on the order of the local velocity dispersion $\sigma$. Defining $B = b/b_{\text{min}} \geq 1$, the encounter rate is then

\[
\frac{d^2\Gamma}{dM_p db} db \sim \frac{dN_p}{dM_p} \frac{d^2\Gamma}{dM_p db} \left(\frac{2GM_p}{v^2}\right)^2 db = G^2 \frac{dN_p}{dM_p} M_p^2 dM_p \frac{2\pi B dB}{v^3}.
\]

The total rate is obtained by integrating over all MP masses and all impact parameters between $b_{\text{min}}$ and $b_{\text{max}}$. Here we are interested in perturbations in the specific angular momentum $J$ of a star relative to the central MBH, and so $b_{\text{max}} \approx r$, the radial distance of the star from the center. MPs with substantially larger impact parameters are much less efficient because their effect on the MBH-star pair is tidal rather than direct.

The relaxation rate due to all MPs at all impact parameters is then

\[
\tau^{-1} = \int_{b_{\text{min}}}^{b_{\text{max}}} db \int dM_p \frac{d^2\Gamma}{dM_p db} \left(\frac{2GM_p}{v^2}\right)^2 db \sim \log \Lambda G^2 \int \frac{dN_p}{dM_p} M_p^2 dM_p,
\]

where $(\delta v/v)^2 \propto 1/B$ is the weighted deflection and $\log \Lambda = \log (b_{\text{max}}/b_{\text{min}})$ is the Coulomb logarithm (here the dependence of $\log \Lambda$ and $v$ on $M_p$ is assumed to be negligible). For stars, typically $\log \Lambda \approx 10$; the omission of large-angle scattering by encounters with $b < b_{\text{min}}$ is thus justified because it introduces only a relatively small logarithmic correction. This formulation of the relaxation time is equivalent to its conventional definition (Spitzer 1987) as the time for a change of order unity in $v^2$ by diffusion in phase space due to scattering, $\tau \sim v^2/D(v^2)$, where $D(v^2)$ is the diffusion coefficient.

If the stars and MPs have distinct mass scales with typical number densities $N_s$ and $N_p$ and rms masses $(M_{s}^2)^{1/2}$ and $(M_{p}^2)^{1/2}$, then MPs dominate if the ratio of the second moments of the MP and star mass distributions, $\mu_2 = N_p (M_{p}^2)/N_s (M_{s}^2)$, satisfies $\mu_2 \gg 1$ (note that for a continuous mass spectrum, this condition is equivalent to $-d \log N/d \log M < 2$).

As discussed in detail in § 3, the central ∼100 pc of the Galactic center contain $10^5 - 10^6 M_\odot$ in stars and about $10^6 - 10^8 M_\odot$ in MPs such as giant molecular clouds or open clusters of masses $10^3 - 10^8 M_\odot$ (Oka et al. 2001; Figer et al. 2002, 2004; Vollmer et al. 2003; Güsten & Philipp 2004; Borissova et al. 2005). An order-of-magnitude estimate indicates that MPs in the GC can reduce the relaxation time by several orders of magnitude:

\[
\frac{\tau_{r,\star}}{\tau_{r,\text{MP}}} = \mu_2 \approx \left(\frac{N_p M_p}{N_s M_s}\right) M_p \approx 10^4 \left(\frac{N_s M_s/N_p M_p}{10}\right)^{-1} M_p/M_s.
\]

Note that $\mu_2$ does not include possible modifications in the value of $\log \Lambda$ for MPs due to their much larger size, which may decrease this ratio by $O(10)$. This estimate is borne out by more detailed calculations (Fig. 1 and Tables 1 and 2), using the formal definition $\tau_r = v^2/D(v^2)$ with $M_s \rho_s \rightarrow \int (dN_p/dM_p)M_p^2 dM_p$.
period, leaving the loss cone always nearly empty. In this empty loss cone regime, the loss cone is relatively large, but the refilling rate is set by the long relaxation timescale (e.g., Lightman & Shapiro 1977):

$$\left( \frac{d\Gamma}{dE} \right)_{\text{empty}} \approx \frac{N_c(E)}{\log (J_c/E) t_r} = \frac{J_D^2(E)}{J_c^2} \frac{1}{\log [J_c(E)/J_c]} \frac{N_c(E)}{P(E)},$$

where $N_c(E)$ is the stellar number density per energy interval.

Far from the MBH (low $E$), $J_D \gg J_c$, stars diffuse into the loss cone many times over one orbit, and the loss cone is always nearly full. In this full loss cone regime, the refilling rate is set by the short orbital time, but the loss cone is relatively small:

$$\left( \frac{d\Gamma}{dE} \right)_{\text{full}} \approx \frac{J_D^2}{J_c^2(E) P(E)}.$$  

Note that here and elsewhere we make the simplifying approximation that the period is a function of energy only, which is true only for motion in a Keplerian potential.

The total contribution to loss cone refilling is dominated by stars with energies near the critical energy $E_c$ (equivalently, the critical typical radius $r_c$), separating the two regimes (Lightman & Shapiro 1977; see § 4). Within $r_c (E > E_c)$, an object, once deflected into the loss cone, can avoid being scattered out of it before reaching the MBH. The empty and full loss cone regimes of infall processes can be interpolated to give a general approximate expression for the differential event rate for these noncoherent encounters (e.g., Young 1977):

$$\frac{d\Gamma}{dE} \approx \frac{J_D^2(E)}{J_c^2} \frac{N_c(E)}{P(E)}, \quad \left( \frac{d\Gamma}{dE} \right)_{\text{full}} \approx \frac{J_D^2}{J_c^2(E) P(E)}.$$  

where $j$ is the loss cone–limited angular momentum change per orbit, which expresses the fact that the loss cone can at most be completely filled during one orbit.

#### 2.1. Noncoherent Loss Cone Refilling

The Fokker-Planck approach to the loss cone problem (e.g., Cohn & Kulsrud 1978) assumes that the effects of multiple small perturbations on the orbit of a test star dominate over the rarer strong close encounters ($b_{\text{max}}/b_{\text{min}} \gg 1$) and that the cumulative effect can be described as diffusion in phase space. The change in the angular momentum of the test star then grows noncoherently, $\Delta J \propto \sqrt{t}$. The change over one orbital period $P$ is $J_D = J_c(E)(P/t_r)^{1/2}$, where $J_c = (2(\psi - E)^{1/2} r)$ is the maximal (circular) angular momentum for a stellar orbit of specific relative energy $E = -r^2/2 + \psi(r)$ and $\psi = -\phi$ is the negative of the gravitational potential, so that $E > 0$ for bound orbits. The magnitude of $J_D$ relative to the $J$ magnitude of the loss cone,

$$J_c \approx \sqrt{2GM*q},$$

where $q$ is the periapse distance to the MBH, determines the mode of loss cone refilling. The relative volume of phase space occupied by the loss cone, $J_D^2/J_c^2(E)$, increases with $E$ (decreases with $r$), while $P$ decreases. Near the MBH (high $E$), $J_D \ll J_c$, stars diffuse slowly into the loss cone and are promptly destroyed over an orbital

![Fig. 1.— Relaxation time as a function of distance from the MBH, for stars (solid line) and for each of the four MP models separately, as listed in Table 2: clusters (dash-dotted lines) and GMCs (dashed lines). The discontinuities are artifacts of the assumed sharp spatial cutoffs on the MP distributions. Two-body stellar processes dominate close to the MBH, where no MPs are observed to exist. However, at larger distances massive clumps (at 1.5 pc < $r$ < 5 pc) and GMCs (at 5 pc < $r$ < 100 pc) are much more important. [See the electronic edition of the Journal for a color version of this figure.]

| MP Type               | $r^*$ (pc) | $N_p$ | $M_p$ ($M_\odot$) | $(M_p^2)^{1/2}$ ($M_\odot$) | $R_p$ (pc) | References       |
|----------------------|------------|-------|-------------------|-----------------------------|------------|------------------|
| Observed GMCs        | <100       | ~100  | $10^4$–$10^6$     | $3 \times 10^6$–$3 \times 10^7$ | 5          | 1, 2             |
| Observed clusters    | <100       | ~10   | $10^2$–$10^5$     | $4.8 \times 10^5$–$2.4 \times 10^4$ | 1          | 3, 4, 5, 6, 7    |
| Observed clumps      | 1.5–3      | ~25   | $10^2$–$10^5$     | $3.7 \times 10^5$–$4.1 \times 10^4$ | 0.25       | 8, 9             |

* Projected distance range enclosing observed MPs.

References.—(1) Oka et al. 2001; (2) Güsten & Philipp 2004; (3) Figer et al. 1999; (4) Figer et al. 2002; (5) Figer 2004; (6) Maillard et al. 2004; (7) Borissova et al. 2005; (8) Genzel et al. 1985; (9) Christopher et al. 2005.
error introduced by neglecting encounters with than the typical value for relaxation by stars. Nevertheless, the fractional contributions from different individual encounters, evaluation lies beyond the scope of this study.

The differential rate is estimated simply by \( \frac{d\Gamma}{db} = N_p v^2 \). When \( P \int_{b_{\text{min}}}^{b_{\text{max}}} db (d\Gamma/db) > 1 \), with \( b_{\text{max}} = r \), all encounters with \( b > b_1 \) are defined as frequent encounters that occur more than once per orbit and add noncoherently.5 Note that even when \( P \int_{b_{\text{min}}}^{b_{\text{max}}} db (d\Gamma/db) > 1 \) for all MPs, perturbations by rare, very massive MPs may still occur less than once per orbit. Our treatment is approximate. A complete statistical treatment of this situation lies beyond the scope of this study.

When the typical number of encounters per orbit is less than 1, the fractional contributions from different individual encounters, \( \delta J \), should be averaged coherently (\( \Delta J \propto \theta^2 \)), subject to the limit that each encounter can at most fill the loss cone. The loss cone–limited change in angular momentum per orbit due to rare encounters is therefore

\[
\frac{\Delta J}{r} = \left[ P \int_{b_{\text{min}}}^{b_{\text{max}}} \frac{d\Gamma}{db} \min(\delta J, J_c) \right]^2 .
\]

In contrast, frequent uncorrelated collisions add up noncoherently (\( \Delta J \propto \sqrt{r} \)), and it is only their final value that is limited by the loss cone (individual steps \( \delta J \) may exceed \( J_c \), but they can then be partially canceled by opposite steps during the same orbit). The loss cone–limited change in angular momentum per orbit is therefore

\[
\frac{\Delta J}{r} = \min \left[ \frac{1}{\log(J_c/J_{\text{lim}})} \int_{b_{\text{min}}}^{b_{\text{max}}} \frac{d\Gamma}{db} \delta J^2 \right] .
\]

The total loss cone–limited angular momentum change per orbit is then approximated by

\[
\frac{\Delta J}{r} = \min \left[ \frac{1}{\log(J_c/J_{\text{lim}})} \int_{b_{\text{min}}}^{b_{\text{max}}} \frac{d\Gamma}{db} \delta J^2 \right] .
\]

and the differential event rate is calculated by equation (7). The contribution of rare encounters is evaluated in the impulse approximation by setting \( \delta J \sim GMpr/r^2 \) in equation (11). We find that the contribution by GC MPs (§3) is generally small. Frequent encounters are the regime usually assumed in the Fokker-Planck treatment of the loss cone problem (e.g., Lightman & Shapiro 1977). To evaluate the contribution of frequent encounters, we do not calculate \( \delta J \) directly, but instead calculate the subexpression \( I = \int P \int_{b_{\text{min}}}^{b_{\text{max}}} db (d\Gamma/db) \delta J^2 \) in equation (12) in terms of the \( b \)-averaged diffusion coefficient \( D(v^2) \), after averaging over the orbit between the periapse \( r_p \) and apoapse \( r_a \) and averaging over the perturber mass function (this is essentially equivalent to the definition of \( D \) in terms of \( t_v \); §2.1),

\[
I = \int dM_p \left[ 2 \int_{v_p}^{r_a} r^2 \Delta v^2 \frac{d\Delta v^2}{v} \right] .
\]

The assumptions involved in the last approximate term (Magorrian & Tremaine 1999) are that the star is on a nearly radial orbit \( (v_p \rightarrow r, v_p \rightarrow 0, r_a \rightarrow 2r) \) and that \( D(v^2) \) (the diffusion coefficient of the transverse velocity relative to the MBH) can be approximated by \( D(\Delta v^2) \) (the diffusion coefficient of the transverse velocity relative to the stellar velocity \( v \)), given explicitly by (Binney & Tremaine 1987, their eq. [8-68])

\[
D(\Delta v^2) = \frac{8\pi G^2 (dN_p/dM_p)M_p^2 \ln \Lambda}{v} K \left( \frac{v}{v^2 \sigma^2} \right) .
\]
where \( K(x) \equiv \text{erf}(x)(1 - 1/2x^2) + \exp(-x^2)/\sqrt{\pi x} \) and where a spatially homogeneous distribution of MPs with a Maxwellian velocity distribution of rms one-dimensional velocity \( \sigma \) was assumed.

To summarize, the event rates are calculated as follows. For each perturber model (Table 2), we integrate over the stellar distribution \( N_s = 1.2 \times 10^4 (r/0.4 \text{ pc})^{-2} \), for \( r > 0.4 \text{ pc} \) and \( M_s = 1 \text{ M}_\odot \) in terms of \( r \), using \( N_s \) to derive the appropriate density of perturbed objects (single stars, \( \S 4.1 \); binaries, \( \S 4.2 \)). At each \( r \), we calculate \( b_{\min} \) (eq. [9]), \( b_1 \) (eq. [10]), \( j_0 \) (eq. [11]), and \( j_E \) (eq. [12]). The integral \( I \) (eq. [13]) is evaluated by taking \( \nu^2 \rightarrow GM(<r)/r \) and correcting approximately for the difference relative to the exact calculation. We use \( b_l \) (eq. [14]) to calculate the differential event rate \( d\Gamma/d\nu \) (eq. [7]), with the \textit{Ansatz} that \( E \rightarrow GM(<r)/2a \), where \( a \equiv (4/5)r \) is an effective semimajor axis, which is motivated by the fact that for a Keplerian isothermal eccentricity distribution, \( (r) = a(1 + (\nu^2)/2) = (5/4)a \). The total event rate is calculated by carrying the integration over \( r \) in the region where \( M_s \) exist, between \( r_{\text{MP}} \) and \( r_{\text{out}} \) (\S 3).

3. MASSIVE PERTURBERS IN THE GALACTIC CENTER

MPs can dominate relaxation only when they are massive enough to compensate for their small space densities. Here we consider only MPs with masses \( M_p \geq 10^5 \text{ M}_\odot \). Such MPs could be molecular clouds of different masses, in particular giant molecular clouds (GMCs), open or globular stellar clusters, and perhaps also IMBHs. As discussed below, observations of the Galaxy reveal enough MPs to dominate relaxation in the central 100 pc. We adopt here a conservative approach and include in our modeling only those MPs that are directly observed in the Galaxy, namely, GMCs and young clusters. We briefly discuss theoretical predictions for two other classes of MPs, dynamically evolved “submerged” clusters and IMBHs, that could well be common in galactic centers and contribute to efficient relaxation.

The dynamically dominant MPs are GMCs. Emission-line surveys of the central \( \sim 100 \text{ pc} \) reveal \( \sim 100 \text{ GMCs} \) with estimated masses in the range \( 10^{4.5} - 5 \times 10^5 \text{ M}_\odot \) and sizes of \( R_p \sim \) a few pc (Miyazaki & Tsuboi 2000; Oka et al. 2001; Güsten & Philipp 2004). We selected for analysis individual, reliably identified GMCs in the central 0.7” of the Galaxy (\( \sim 100 \text{ pc} \) of the GC) from the sample observed by Oka et al. (2001). Figure 2 shows the empirical GMC mass function following the Oka et al. (2001) virial mass estimates as an upper mass limit and adopting a lower limit 10 times smaller, following Miyazaki & Tsuboi (2000), who found that LTE mass estimates are typically an order of magnitude lower than the virial ones. Note that the more recent GMC CO(1–0) molecular line observations by Oka et al. (2001), which we use here, indicate a more massive and flatter mass function than that derived for their earlier CS(1–0) molecular line observations (Miyazaki & Tsuboi 2000). This is probably due to the higher sensitivity of the CO(1–0) line to lower density molecular gas (M. Tsuboi 2006, private communication).

Inside the inner 5 pc of the GC (a volume smaller or comparable to that of a GMC or a stellar cluster), the most massive local structures are the molecular gas clumps observed in the circumnuclear gaseous disk and its associated spiral-like structures (Genzel et al. 1985; Christopher et al. 2005). Their sizes are \( \sim 0.25 \text{ pc} \) and their masses are estimated to be in the range \( 10^5 - 10^6 \text{ M}_\odot \), where the lower estimates are based on the assumption of optically thin HCN(1–0) line emission and the upper estimates are based on the optically thick assumption, which also coincides with the virial estimates.

It is possible to obtain a model-independent estimate of the effect of the \textit{observed} MPs on the relaxation time by directly calculating the quantity \( \mu_{\text{obs}} = \sum M_p^2/nM_p^2 \), which is listed in Table 2. The observed GMC masses show that \( \mu_{\text{obs}} \sim 2 \times 10^6 - 10^8 \) on the 100 pc scale and \( \mu_{\text{obs}} \sim 6 \times 10^1 - 10^3 \) on the 5 pc scale. This clearly indicates that MPs dominate the relaxation on all relevant length scales where MPs exist. For the purpose of our numeric calculations below, it is convenient to describe the differential mass function analytically.

Here we model the data as either a power-law distribution (d\( n_p/dM_p \propto M_p^{-\beta} \)) or a lognormal (LN) distribution (d\( n_p/dM_p \propto \exp[-(\log M_p - \mu_S)/2\sigma_S^2]/(2\pi)^{1/2}M_p\sigma_S \), where \( \mu_S \) and \( \sigma_S \) are the mean and standard deviation of \( \log M_p \)), whichever fits the data better. Figure 2 shows that the power-law distribution is a good fit for the mass function of the GMCs (Miyazaki & Tsuboi 2000) and for the lower estimate of the gas clumps. For the upper estimate of the molecular gas clumps, the mass function is better fitted by a LN distribution. It should be emphasized that our results and conclusions are determined primarily by the large values of \( \mu_S \), not by the detailed form of the mass function. We repeated our calculations with several alternative distributions and found qualitatively similar results. The assumed high-mass cutoff of the mass function is important, as it determines the magnitude of \( \mu_S \). Figure 2 shows that our models do not extrapolate beyond the maximum observed MP masses (and even fall below them for GMCs). Our analytical models thus provide a faithful and robust representation of the MP mass function.

We obtained best-fit power-law indices \( \beta \) and best-fit LN moments \( \mu_S \) and \( \sigma_S \) for the lower and upper mass estimates and then chose the mass function model that best described the data. The cumulative mass functions and best fits are shown in Figure 2 and are listed in Table 1. Both the GMCs’ upper mass estimate (1.4 \( \times 10^5 \text{ M}_\odot \leq M_p \leq 5 \times 10^7 \text{ M}_\odot \)) and the lower mass estimates (1.4 \( \times 10^4 \text{ M}_\odot \leq M_p \leq 5 \times 10^6 \text{ M}_\odot \)) are well fitted by a power-law distribution with \( \beta = 1.2 \) (models GMC1 and GMC2 in Table 2). The clumps’ lower mass estimates (2.4 \( \times 10^2 \text{ M}_\odot \leq M_p \leq 1.1 \times 10^4 \text{ M}_\odot \)) are best fitted by a power-law distribution with \( \beta \approx 1.1 \). The clumps’ upper mass estimates
stars with a number density distribution \( N_*(r) \propto r^{-2} \) outside the inner 1.5 pc (i.e., a constant MP-to-star ratio). (4) Mutually exclusive perturber types (i.e., a single type of perturber is assumed to dominate relaxation, as indicated by the detailed calculations presented in Fig. 1).

The five perturber models are detailed in Table 2: Stars, Clusters1, Clusters2, GMC1, and GMC2, which represent, respectively, the case of relaxation by stars only, by heavy or light stellar clusters, and by heavy GMCs or light GMCs. Table 2 also lists \( \mu_\text{obs} \), the observed ratio of the second moment of the various MPs to that of the stars.

4. MASSIVE PERTURBER–DRIVEN INTERACTIONS WITH A MBH

The maximal differential loss cone refilling rate, which is also the close encounter event rate, \( d\Gamma/dE \), is reached when relaxation is efficient enough to completely refill the loss cone during one orbit (eq. [7]). Further decrease in the relaxation time does not affect the event rate at that energy. MPs can therefore increase the differential event rate over that predicted by stellar relaxation only at high enough energies, \( E > E_c \) (equivalently, at small enough typical radii, \( r < r_c \)), where slow stellar relaxation fails to refill the empty loss cone. The extent of the empty loss cone region increases with the maximal periapse \( q \), which in turn depends on the close encounter process of interest. For example, the tidal disruption of an object of mass \( M \) and size \( R \) occurs when \( q < r_t \), the tidal disruption radius:

\[
r_t \approx R(M_*/M)_{1/3}.
\]

This approximate disruption criterion applies both for single stars \((M = M_*, R = R_*)\) and for binaries, where \( M \) is the combined mass of the binary components and \( R \) is the binary’s semimajor axis, \( a \). Stellar radii are usually much smaller than typical binary separations, but stellar masses are only \( \approx 2 \) times smaller than binary masses. Binaries are therefore disrupted on larger scales than single stars. In the GC this translates to an empty (stellar relaxation) loss cone region extending out to \( r^* < 3 \) pc for single stars and out to \( r^b > 100 \) pc for binaries. In the GC, \( r_{\text{MP}} \ll r^* \ll r^b \), and so MPs are expected to increase the binary disruption rate by orders of magnitude, but increase the single-star disruption rate only by a small factor. This is depicted qualitatively in Figure 3, which shows the local rates \((d\Gamma/d\log r)\) for disruption of single stars and binaries due to stellar relaxation or relaxation by a simplified one-component MP model.

Figure 3 also shows that the MP-induced disruption rate is dominated by binaries originating near the inner cutoff, \( r_{\text{MP}} \) [in the following discussion the initial orbital energy of the disrupted objects is expressed by the corresponding radius of a circular orbit, \( r \sim GM(<r)/2E \)]. This is qualitatively different from the usual case of stellar tidal disruption induced by stellar relaxation, which mainly occurs inside \( r^* \) and is dominated by stars originating near the outer limit, which is \( \min(r^*, r_b) \); usually \( r^* \approx r_b \) (Lightman & Shapiro 1977). The difference can be understood by considering the \( r \) dependence of \( d\Gamma/d\log r \). Neglecting logarithmic terms, the empty and full local loss cone rates are, respectively (eqs. [5] and [6]),

\[
\frac{d\Gamma_e}{d\log r} \sim \frac{N_c(<r)}{r(r)},
\]

\[
\frac{d\Gamma_f}{d\log r} \sim \frac{M_*}{M_* + M_c(<r)} \frac{1}{r} \frac{q N_*(<r)}{r^2},
\]

where \( N_*(<r) \) is the number of stars enclosed within \( r \).
The empty loss cone regime for binary-MBH interactions extends out to \( r > 100 \) pc because of their large tidal radius. On these large scales MPs are abundant enough to dominate the relaxation processes. Here we focus on three-body exchange interactions (Hills 1991, 1992; Yu & Tremaine 2003), which lead to the disruption of the binary, the energetic ejection of one star, and the capture of the other star on a close orbit around the MBH.

Various phenomena associated with such exchange interactions were suggested and explored, Hills (1988), and later Yu & Tremaine (2003), Gualandris et al. (2005), Ginsburg & Loeb (2006), and Bromley et al. (2006), studied hypervelocity stars ejected from the GC following tidal disruption by the MBH. Gould & Quillen (2003) suggested this mechanism to explain the origin of the young stars near the Galactic MBH. Miller et al. (2005) proposed that compact objects captured following a binary disruption event will eventually be sources of GWs from zero-eccentricity orbits, in contrast to high-eccentricity sources typical of single-star inspiral (Hopman & Alexander 2005).

The event rates estimated by these authors vary substantially. Hills (1988) assumed a full loss cone and a fraction \( f_{\text{bin}} = 0.02 \) of the stars in binaries with small enough semimajor axes to produce a high-velocity star (\( a < 0.1 \) AU) and derived a three-body exchange rate of \( \sim 10^{-3} f_{\text{bin}} \) yr\(^{-1}\). Yu & Tremaine (2003) took into account the empty loss cone regime and argued for a higher fraction of relevant binaries (\( f_{\text{bin}} = 0.04 \) for binaries with \( a < 0.3 \) AU that can survive 0.8 pc from the MBH), thereby obtaining a rate of \( \sim 2.5 \times 10^{-3} f_{\text{bin}} \) yr\(^{-1}\), 3 orders of magnitude smaller than that estimated by Hills. These calculations assumed the same binary separation for all binaries and a constant binary fraction at all distances from the MBH (two possibilities were considered, \( a = 0.3 \) AU and \( a = 0.03 \) AU).

The binary fraction and typical binary semimajor axis depend on the binary mass and on the rate at which binaries evaporate by encounters with other stars. This depends in turn on the stellar densities and velocities and therefore on the distance from the MBH. Here we take these factors into account and estimate in detail the three-body exchange rate for MP-driven relaxation. The rate is proportional to the binary fraction in the population, which is the product of the poorly known binary initial mass function (IMF) in the GC and the survival probability against binary evaporation.

We assume for simplicity equal-mass binaries, \( M_{\text{bin}} = 2M_* \), since the observations indicate that moderate mass ratios dominate the binary population (Duquennoy & Mayor 1991; Kobulnicky et al. 2006). The evaporation timescale at a distance \( r \) from the center for a binary of semimajor axis \( a \) composed of two equal-mass stars of mass \( M_* \) and lifetime \( t_* \) is (e.g., Binney & Tremaine 1987)
\[
t_{\text{evap}}(M_*, a, r) = \frac{M_{\text{bin}}}{(M_*) 16 \sqrt{\pi} \rho(r) G a \ln \Lambda_{\text{bin}}}.
\]

The Coulomb factor for binary evaporation, \( \Lambda_{\text{bin}} = a^2/4G(M_*) \), expresses the fact that the binary is only affected by close perturbations at distances smaller than \( a/2 \). The MPs considered...
here are extended objects (Table 2) and therefore do not affect the binary evaporation timescale (IMBH MPs could be a possible exception). Although binary evaporation is a stochastic process and the actual time to evaporation differs from binary to binary, we expect a small scatter in the evaporation rate, \( \Delta t_{\text{evap}}/t_{\text{evap}} \ll 1 \), because evaporation is a gradual process caused by numerous weak encounters. Evaporation is thus better approximated as a fixed limit on the binary lifetime rather than as a Poisson process (where \( df/dt = \frac{1}{t_{\text{bin}}} \)).

The present-day binary semimajor axis distribution is therefore

\[
d\frac{N_{\text{bin}}}{da} \bigg|_{r_{\text{th}}} = f_{\text{bin}}(M_*) \frac{df}{da} \bigg|_{r_{\text{th}}} N_* \min \left( \frac{t_{\text{evap}}(M_*, a, r)}{\min(t_{\text{th}}, t_*)} \right),
\]

where \( a \) is the semimajor axis of the infalling binary and \( a_1 \) is that of the captured star (the MBH-star "binary"). Most values of \( a_1 \) fall within a factor of 2 of the mean. This relation maps the semimajor axis distribution of the infalling binaries to that of the captured stars: the harder the binaries, the more tightly bound the captured stars. The velocity at infinity of the ejected star (neglecting the Galactic potential) is

\[
v_{\infty} = 2^{1/2} \left( \frac{GM_{\text{bin}}}{a} \right) \left( \frac{M_*}{M_{\text{bin}}} \right)^{1/3} \propto M_{\text{bin}}^{2/3}/a \quad \text{(an equal-mass binary with periapse at } r_1 \text{ is assumed; Hills 1988).}
\]

The harder the binary, the higher the value of \( v_{\infty} \). The periapse of the captured star is at \( r_1 \), and therefore its eccentricity is very high (Hills 1991, 1992; Miller et al. 2005): \( e = 1 - r_1/a_1 \geq 1 - 1.8(M_{\text{bin}}/M_*)^{1/3} \geq 0.98 \) for values typical of the GC.

We now consider separately the implications of three-body exchange interactions of the MBH with old \((t \gtrsim t_{\text{th}})\) binaries and massive young \((t < 5 \times 10^7 \text{ yr})\) binaries.

### 4.2.1. Low-Mass Binaries

The properties of binaries in the inner GC are at present poorly determined. The period distribution of solar neighborhood binaries can be approximated by a lognormal distribution with a median period of 180 yr \((a \sim 40 \text{ AU; Duquennoy \\& Mayor 1991})\). The total binary fraction of these binaries is estimated at \( f_{\text{bin}} \sim 0.3 \) (Lada 2006). Adopting these values for the GC, the total binary disruption rate by the MBH can then be calculated, as described in §2, by integrating \( dN_{\text{bin}}/da \) (eq. [20]) over the binary \( a \)-distribution and over the power-law stellar density distribution of the GC up to 100 pc (Genzel et al. 2003). Table 3 lists the capture rates for the different perturber models, assuming a typical old equal-mass binary of \( M_{\text{bin}} = 2 M_0 \).

The old, low-mass binary disruption rate that we derive for stellar relaxation alone is \( \sim 5 \times 10^{-7} \text{ yr}^{-1} \), which is \( \sim 5 \) times lower than, but still in broad agreement with, the result of Yu \\& Tremaine (2003). Their rate is somewhat higher because they assumed a constant binary fraction and a constant semimajor axis for all binaries, even close to the MBH, where these assumptions no longer hold.

MPs increase the binary disruption and high-velocity star ejection rates by factors of \( \sim 10^3 - 10^5 \) and effectively accelerate stellar migration to the center. This can modify the stellar distribution close to the MBH by introducing a "source term" to the stellar current into the MBH. Low-mass stars are at present too faint to be directly observed in the GC. However, such a source term may have observable consequences, since it can increase the event rate of single-star processes such as tidal disruption, tidal heating, and GW emission from compact objects, in particular from compact objects on zero-eccentricity orbits (Miller et al. 2005; in contrast, GW from inspiraling single stars occurs on high-eccentricity orbits; Hopman \\& Alexander 2005). We calculated numerical solutions of the Fokker-Planck equation for the stellar distribution around the MBH with a captured-star source term. These preliminary investigations (H. B. Perets et al. 2007, in preparation) confirm that the total accumulated mass of captured stars does not exceed the dynamical constraints on the
extended mass around the MBH (Mouawad et al. 2005), because two-body relaxation and likely also resonant relaxation (Rauch & Tremaine 1996; Hopman & Alexander 2006a) scatter enough of them into the MBH or to wider orbits.

4.2.2. Young Massive Binaries

MPs may be implicated in the puzzling presence of a cluster of main-sequence B stars \(4 M_\odot \lesssim M_\star \lesssim 15 M_\odot\) in the inner \(\sim 1''\) (\(\sim 0.04\) pc) of the GC. This so-called S cluster is spatially, kinematically, and spectroscopically distinct from the young, more massive stars observed farther out, on the \(\sim 0.05\) – \(0.5\) pc scale, which are thought to have formed from the gravitational fragmentation of one or two gas disks (Levin & Beloborodov 2003; Genzel et al. 2003; Milosavljević & Loeb 2004; Paumard et al. 2006). There is, however, still no satisfactory explanation for the existence of the seemingly normal young massive main-sequence stars of the S cluster, so close to a MBH (see review of proposed models by Alexander 2005; also a recent model by Levin 2007).

Here we revisit an idea proposed by Gould & Quillen (2003), which is that the S stars were captured near the MBH by three-body exchange interactions with infalling massive binaries. Originally, this exchange scenario lacked a plausible source for the massive binaries. Gould & Quillen speculated that they originated in an unusually dense and massive young cluster on an almost radial infall trajectory, but they concluded that such a finely tuned scenario seems unlikely. Furthermore, a massive cluster is expected to leave a tidally stripped tail of massive stars beyond the central \(0.5\) pc (Kim et al. 2004; Gürkan & Rasio 2005), which are not observed (Paumard et al. 2006). Alternatively, it must contain an unusually massive central MBH to hold it together against the tidal field of the GC (Hansen & Milosavljević 2003). However, such a massive MBH is well beyond what is predicted by simulations of IMBH formation by runaway collisions (Gürkan et al. 2004; Gürkan & Rasio 2005) or anticipated by extrapolating the \(M-\sigma\) relation (Ferrarese & Merritt 2000; Gebhardt et al. 2000) to clusters.

MP-driven three-body exchanges circumvent the problems of the cluster infall scenario by directly bringing massive field binaries to \(r_\star\), without requiring massive clusters of unusual, perhaps even impossible properties. The ongoing star formation in the central \(\sim 100\) pc implies the presence of a large reservoir of massive stars there, which are indeed observed in the central few tens of pc both in dense clusters and in the field (Figer et al. 1999; Figer 2003; Munu et al. 2006). It is plausible that a high fraction of them are in binaries. It is assumed here that the field binaries are dynamically relaxed, which is reasonable in view of the fact that MP relaxation times are shorter or comparable to their typical stellar lifetimes over most of the range of GMC-dominated MP models.

We model the binary population of the GC in the S star mass range, \(4 M_\odot \lesssim M_\star \lesssim 15 M_\odot\), by assuming equal-mass binaries that follow the single-star mass function with an initial binary fraction of \(f_{bin} \sim 0.75\), as observed elsewhere in the Galaxy (Lada 2006; Kobulnicky et al. 2006). Because the stellar evolutionary life span of such stars is relatively short, massive binaries are essentially unaffected by dynamical evaporation. We assume star formation at a constant rate for 10 Gyr with a Miller-Scalo IMF (Miller & Scalo 1979) and use a stellar population synthesis code (Stenberg et al. 2003) with the Geneva stellar evolution tracks (Schaller et al. 1992) to estimate that the present-day number fraction of stars in the S star mass range is \(3.5 \times 10^{-4}\) (and less than 0.01 of that for \(M_\star > 15 M_\odot\) stars). Note that if star formation in the GC is biased toward massive stars (Figer 2003; Stolte et al. 2005), this estimate should be revised upward. We adopt the observed solar neighborhood distribution of the semimajor axis of massive binaries, which is lognormal with \(\langle \log a \rangle = -0.7 \pm 0.6\) AU (i.e., 63% of the binaries with \(a = 0.2^{+0.15}_{-0.10}\) AU; 91% with \(a = 0.2^{+0.26}_{-0.19}\) AU; Garmany et al. 1980; Kobulnicky et al. 2006). Massive binaries are thus typically harder than low-mass binaries and will be tidally disrupted (eq. [16]) closer to the MBH and leave a more tightly bound captured star.

We represent the massive binaries by binaries composed of equal-mass stars in the midrange of the S star masses, with \(M_{bin} = 15 M_\odot\) and \(t_\star (7.5 M_\odot) \sim 5 \times 10^7\) yr, and integrate over the stellar distribution and the binary \(a\)-distribution as before, to obtain the rate of binary disruptions, \(\Gamma\), the mean number of captured massive stars in a steady state, \(N_{\star,c} = \Gamma t_\star\), and their semimajor axis distribution (eq. [21]). Table 3 compares the number of captured young stars in a steady state for the different MP models, both on the \(r < 0.04\) pc scale (the S cluster) and on the \(0.04 < r < 0.4\) pc scale (the stellar rings), with current observations (Eisenhauer et al. 2005; Paumard et al. 2006).

The S stars are found in the central \(<0.04\) pc (Ghez et al. 2005; Eisenhauer et al. 2005). If they were captured by binary disruptions, they must have originated from massive binaries with \(a \lesssim 3.5\) AU. This is consistent with the semimajor axis distribution of massive binaries. The number of captured massive stars falls rapidly beyond 0.04 pc (Table 3) because wide massive binaries are rare. This capture model thus provides a natural explanation for the central concentration of the S cluster (Fig. 4). The absence of more massive stars in the S cluster \((M_\star > 15 M_\odot\); spectral type O V) is a statistical reflection of their much smaller fraction in the binary population. Figure 4 and Table 3 compare the cumulative semimajor axis distribution of captured B stars, as predicted by the different MP models, with the total number of young stars observed in the inner 0.04 pc (\(\sim 35\) stars; Eisenhauer et al. 2005; Ghez et al. 2005; Paumard et al. 2006). Of these, only \(\sim 10\) have full orbital solutions (in particular, solutions for \(a\) and \(e\)) at present. For the others, we assume the Ansatz that \(a\) is similar to the observed projected position. The numbers predicted by the GMC-dominated MP models are consistent with the observations, unlike the stellar relaxation model, which falls short by 2 orders of magnitude.
The binary capture model predicts that captured stars have very high initial eccentricities. Most of the solved S star orbits do have $e > 0.9$, but a couple have $e \sim 0.3-0.4$ (Eisenhauer et al. 2005). Normal, noncoherent stellar relaxation is slow, even after taking into account the decrease in $r_i$ toward the center due to mass segregation (Hopman & Alexander 2006b). It is unlikely that it could have decreased the eccentricity of these stars over their relatively short lifetimes. However, the much faster process of resonant relaxation (Rauch & Tremaine 1996) may be efficient enough to randomize the eccentricity of a fraction of the stars, and it could thus possibly explain the much larger observed spread in eccentricities (Hopman & Alexander 2006a). Additional orbital solutions and a better estimate of the efficiency of resonant relaxation in the GC are required for more detailed comparisons between observations and the MP model predictions.

### 4.2.3. Hypervelocity Stars

Each captured star is associated with an ejected companion, which in some cases is launched with a very high velocity. The one-to-one correspondence between the number of captured S stars and the number of early-type hypervelocity stars (HVSs) is thus a generic prediction of binary capture models. The MP capture scenario specifically implies the continuous and isotropic ejection of both young and old HVSs from the GC. Recent observations of HVSs (Brown et al. 2005; Hirsch et al. 2005; Edelmann et al. 2005; Brown et al. 2006a, 2006b) are consistent with a GC origin and favor a steady state temporal distribution and an isotropic spatial distribution over a burstlike nonspherical distribution of these HVSs by equal-mass binaries of mass $\sim 10^{-8}$ (Eisenhauer et al. 2005; Brown et al. 2006a). We model the parent binaries of these HVSs by equal-mass binaries of mass $\sim 8 M_\odot$ and $t_e(4M_\odot) = 2 \times 10^8$ yr. The ejection velocity was found in numerical simulations (Hills 1988; Bromley et al. 2006) to scale as

$$v_\infty = 1776 \text{ km s}^{-1} \left( \frac{a}{0.1 \text{ AU}} \right)^{-1/2} \left( \frac{M_{\text{bin}}}{2 M_\odot} \right) \left( \frac{M_\odot}{3.7 \times 10^6 M_\odot} \right)^{1/6}. $$

(22)

To reproduce the high HVS velocities, we consider binaries with $a < 1$ AU, which are tidally disrupted at $r_i < 3.7 \times 10^{-4}$ pc and eject a HVS with $v_\infty \approx 900$ km s$^{-1}$, the escape velocity from the bulge (Haardt et al. 2006). Taking only the GCMCs into account, we predict that tens to hundreds of such HVSs exist in the Galaxy, in agreement with the deduced HVS populations, whereas stellar relaxation predicts only 1.3 such stars (see Table 3).

We note that Brown et al. (2006a) used the $10^{-5}$ yr$^{-1}$ total rate of hypervelocity star ejection calculated by Yu & Tremaine (2003; including binaries of all stellar types, assuming stellar relaxation only, and normalized to a fiducial 10% binary fraction) to estimate the number of late B V HVSs in a Salpeter IMF at 10–25 $M_\odot$. This theoretical prediction seems to be in rough agreement with the observations (and contradicts our much lower estimate of 1.3 HVSs). However, the rates of Yu & Tremaine are inapplicable here and lead to a significant overestimate of the number of HVSs because their binary population model is not appropriate for massive binaries in the GC. On the one hand, the young binary population does not extend all the way to the center, as assumed by Yu & Tremaine for the general binary population (following Paumard et al. [2006], who do not find any B V stars between 0.5 and 1 pc of the GC, we truncate the massive binary population inside 1.5 pc). The massive binary population in the GC is a few times 10 smaller than that implied by a simple scaling of the Yu & Tremaine general binary population. On the other hand, the binary fraction of young massive binaries is 70%, rather than 10% (§ 4.2.2). We conclude that the agreement found by Brown et al. is accidental and that binary disruption by stellar relaxation is insufficient to explain the number of observed HVSs, whereas MP-induced relaxation can reproduce the observations.

### 5. SUMMARY

Relaxation by MPs dominates relaxation by two-body stellar interactions when the ratio between the second moments of their respective mass functions satisfies $\mu_2 = N_p(M_{P}^2)/N_p(M_{P}^2) \gg 1$. We show that Galactic MPs (mostly GMCs and smaller molecular gas clumps that exist outside the inner few pc) dominate and accelerate relaxation in the inner $\sim 100$ pc of the GC. This is plausibly the case in the centers of late-type galaxies in general. There is also evidence for molecular gas in the centers of early-type galaxies (e.g., Rupen 1997; Knapp 1999), which suggests that MPs may dominate relaxation there as well and lead to the relaxation of the central regions of galactic bulges in general.

Relaxation determines the rate at which stars and binaries are deflected to nearly radial (loss cone) orbits that bring them closer to the MBH than some critical periapse $q$, where they undergo a strong destructive interaction with it. The size of $q$ depends on the nature of the interaction of interest (e.g., tidal disruption, three-body exchange). It is much larger for binaries than for single stars due to the binary’s larger effective size.

We extend loss cone theory to approximately treat rare encounters with MPs and apply it to explore the implications of MPs on the rates of different types of close encounters. The rate reaches its maximum when loss cone orbits are replenished by scattering within an orbital time (the full loss cone regime). This is more easily achieved when the phase-space volume of the loss cone is small; that is, when $q$ is small. MPs thus affect only those processes with large $q$ whose loss cones are too large to be efficiently replenished by stellar encounters (the empty loss cone regime).

We show that MPs will not contribute much to the disruption of single stars in the GC, whose loss cone is efficiently replenished by stars outside the central $\sim 2$ pc (MPs may accelerate the consumption of stars by more massive MBHs, where $q$ is significantly larger, or the capture of stars in accretion disks). However, MPs will enhance by factors of 10–1000 the tidal disruption rate of infalling binaries, which result in the capture of one of the stars on a tight orbit around the MBH and the ejection at high velocity of the other star (Hills 1991, 1992; Yu & Tremaine 2003). The enhancement of the event rates is dominated by the innermost MPs, and so the uncertainty in determining the MP distribution on the smallest scales dominates the uncertainties in the total event rate. Detailed observations of MPs in the inner GC allow us to robustly predict their effects in the Galaxy. We show that MP-induced disruptions of relatively rare massive binaries can naturally explain the puzzling presence of apparently normal main-sequence B stars in the central 0.04 pc of the GC (Eisenhauer et al. 2005) and at the same time can account for the observed HVSs that are well on their way out of the Galaxy (Brown et al. 2005; Edelmann et al. 2005; Hirsch et al. 2005; Brown et al. 2006a, 2006b; Bromley et al. 2006). Tidal disruptions of the many more faint low-mass binaries can efficiently supply single stars on very eccentric tight
orbits near the MBH. Such an increase in the number of stars in tight orbit near the MBH may increase the rates of single-star processes such as tidal disruption and heating or GW emission from tightly bound compact objects (Miller et al. 2005).

Finally, MP-driven interactions also have cosmological implications for the coalescence of binary MBHs following galactic mergers (Zhao et al. 2002). We suggest that MPs can accelerate the dynamical decay of binary MBHs by efficiently supplying stars for the slingshot mechanism and thereby help solve the "last parsec" stalling problem. MP-driven loss cone refilling will operate even in the case of a spherical potential, where other suggested mechanisms are inefficient, thus allowing MBHs to coalesce on the dynamical timescale of the galactic merger. A detailed treatment of this idea will be presented elsewhere (H. B. Perets & T. Alexander 2007, in preparation).

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ERRATUM: “MASSIVE PERTURBER–DRIVEN INTERACTIONS BETWEEN STARS 
AND A MASSIVE BLACK HOLE” (ApJ, 656, 709 [2007])

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There is an error in the plotted cumulative mass function of the giant molecular clouds in Figure 2 (top panel, two rightmost curves). Due to a misinterpretation of the mass estimates in Table 1 of T. Oka et al. (ApJ, 562, 348 [2001]), Figure 2 includes 12 (out of 60) data points that are multiple entries for the same clouds, expressing alternative mass estimation assumptions (we thank R. O’Leary for pointing this out and T. Oka for clarifying the issue). The corrected observed ratio of second moments of the mass function $\mu_2$ in Table 2 should be modified downward to $10^8$ for model GMC1 and to $10^6$ for model GMC2. However, the analytic power-law fit for the mass function that was actually used in the numerical calculations was very conservative to begin with, and still conservatively underestimates $\mu_2$ by a factor of 2. Therefore, the results presented in the paper remain unchanged.