Gravitational radiation reaction and inspiral waveforms in the adiabatic limit

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We describe progress evolving an important limit of binary orbits in general relativity, that of a stellar mass compact object gradually spiraling into a much larger, massive black hole. These systems are of great interest for gravitational wave observations. We have developed tools to compute for the first time the radiated fluxes of energy and angular momentum, as well as instantaneous snapshot waveforms, for generic geodesic orbits. For special classes of orbits, we compute the orbital evolution and waveforms for the complete inspiral by imposing global conservation of energy and angular momentum. For fully generic orbits, inspirals and waveforms can be obtained by augmenting our approach with a prescription for the self force in the adiabatic limit derived by Mino. The resulting waveforms should be sufficiently accurate to be used in future gravitational-wave searches.

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The late dynamics of a merging compact binary remains one of the greatest challenges of general relativity (GR). GR doesn’t have a “two body” problem so much as it has a “one spacetime” problem: one must use the Einstein field equations to find the dynamical spacetime describing a multibody system. Although numerical relativity has made great progress in recent years (e.g., [1] and references therein), most astrophysically relevant progress has come from identifying a small parameter which defines a perturbative expansion. One such approach is post-Newtonian (PN) theory (e.g., [2] and references therein) — essentially, an expansion in interaction potential \( \frac{GM}{r} \). PN theory works very well when the bodies’ separation \( r \) is large. As the bodies come close, the expansion must be iterated to high order (though it may be possible to modify the expansion to improve its convergence [3]).

Strong field binaries can be modeled very accurately if one member is far more massive than the other. The spacetime is then that of the larger body plus a perturbation due to the smaller body, with the mass ratio acting as an expansion parameter. This limit is not just of formal interest; binaries consisting of “small” compact bodies (white dwarfs, neutron stars, or black holes with mass \( 1 \ 100 M ) captured onto highly eccentric orbits of massive black holes \( M \ 10^5 \ 10^6 M \) are important targets for space-based gravitational-wave (GW) antennae, especially the LISA mission. Such captures are estimated to occur with a rate density \( 10^4 \) Gpc \(^{-3} \) yr \(^{-1} \). Convolving with LISA’s planned sensitivity, one expects to measure dozens to thousands of events per year [5]. Precisely measuring GW phase as the smaller body spirals into the black hole (driven by GW backreaction) and fitting to detailed models will determine system parameters with extraordinary accuracy. For example, the hole’s mass and spin should be determined to \( < 0.1\% \) [6]. It should even be possible to “map” the black hole’s spacetime, testing whether it satisfies GR’s stringent requirements [7,8].

Thanks to their extremal mass ratio, a formal prescription for modeling such systems now exists. At lowest order, the small body follows a geodesic orbit of the black hole [e.g., Ref. [9], Eqs. (33.32a)–(33.32d)]. This must be corrected by the small body’s interaction with its own spacetime distortion. The electromagnetic manifestation of this self force was given by Dirac [10], yielding the Abraham-Lorentz-Dirac equation of motion. Equations for a body’s curved spacetime self-interaction were worked out by Mino, Sasaki, and Tanaka [11] and by Quinn and Wald [12]; see Ref. [13] for an overview. Many [14] researchers are now working to develop practical schemes to compute the self force in black hole spacetimes, which is quite a challenge.

Adiabatic radiation reaction (ARR). Astrophysical extreme-mass-ratio binaries allow a significant simplification for most of the inspiral, up to the last few orbits before the final plunge and merger. In this regime, the system evolves adiabatically: the radiation reaction time \( T_{\text{rad}} \) is much larger than the orbital time \( T_{\text{orb}} \): \( T_{\text{orb}} / T_{\text{rad}} \approx \mathcal{M} - 1 \). The inspiral can be approximated as a flow through a sequence of geodesic orbits. To compute the leading-order, adiabatic waveforms, it is necessary only to know the time-averaged rates of change of the three constants of motion: the energy \( E \), axial angular momentum \( L_z \), and Carter constant \( Q \) [15,16].

This approximation can be understood by expanding the small body’s self acceleration \( a \). Define \( a = a_{\text{cons}} + a_{\text{diss}} \). We can then write \( a = a_{\text{cons}} + a_{\text{diss}} + a_{\text{cons}} + \cdots \). Terms labeled “diss” describe dissipative aspects of the self force; they drive the inspiral. Those labeled “cons” are conservative, representing non-dissipative components that contribute to the inspiraling body’s inertia. Dissipative terms accumulate secularly; conservative pieces do not. Their observable impact can be seen in the following expression obtained from a two-time expansion [15,16] for the azimuthal orbital phase \( \psi(t) \) (and correspondingly, the GW phase of each harmonic component of the waveform):

\[
\psi(t) = \frac{1}{\pi} \left( \psi_0(t) + \sum_{n=1}^\infty (\psi_n(t) + O(m^2)) \right)
\]

where the leading-order, adiabatic waveforms contain only the term \( n = 0 \) and omit the subleading term \( n = 1 \) which contributes a phase correction of order unity over the entire inspiral. The leading-order term is determined by \( a_{\text{diss}} \). Because it does not accu-
mulate secularly, \(a_0^{\text{cons}}\) does not contribute at leading order; it (along with \(a_0^{\text{vis}}\)) contributes to the subleading term \(4(t)\).

Estimates based on post-Newtonian theory suggest that the leading-order, adiabatic waveforms will be sufficiently accurate to detect waves in LISA data, and to provide an initial estimate of source parameters \(17\). The phase mismatch between \(0(t)\) and the true signal is compensated for by adjustments in model parameters, which introduces systematic error in the inferred parameters. Eliminating this systematic error will require "measurement templates" that accurately model \(4(t)\) — a far more difficult task.

In this paper, we describe recent progress in computing adiabatic waveforms, using (i) an approach called "Poor man’s radiation reaction” (PMRR) for special classes of orbits, and (ii) augmenting PMRR using recent results of Mino \(19\) for fully generic orbits.

**Poor man’s radiation reaction.** PMRR only requires knowledge of the energy and angular momentum the binary radiates; imposing global conservation, we evolve those quantities, approximating inspiral as a flow through a sequence of orbits. PMRR cannot rigorously evolve all the constants \(E, L_z, Q\) describing black hole orbits, except in special cases. It is simple to evolve \(E\) and \(L_z\) — one computes the rate at which \(E\) and \(L_z\) are carried to infinity \(20\) and absorbed by the hole's event horizon \(21\), and imposes global conservation \(22\): \(E_{\text{out}} + E_{\text{rad}} = 0, L_{z\text{out}} + L_{z\text{rad}} = 0\). Carter’s constant \(Q\) cannot be so evolved, since there is no notion of \(Q\)-flux carried by radiation and no associated global conservation law. Orbits with zero eccentricity or inclination are sufficiently constrained that \(Q\) is fixed by \(E\) and \(L_z\) \(22\).

The general case admits no such constraints. Nevertheless we can produce inspirals accurate enough for data-analysis algorithm development using simple, crude approximations based on limiting cases. For example, forcing a certain inclination angle to be constant determines the inspiral using only \(E\) and \(L_z\) fluxes \(24\). Such inspirals are very unlikely to be accurate enough for detection templates.

The central engine of PMRR is a complex function \(4\) representing the small body’s perturbation to the black hole’s curvature. This function is the Weyl (vacuum) curvature tensor \(\mathcal{C}_{abcd}\) projected onto a set of null vectors that are very convenient for describing outgoing radiation. Far away, \(4\) is simply related to the two GW polarizations:

\[
4(t \to 1) = \frac{1}{2} \frac{\partial^2}{\partial t^2} \Theta_{ll} = 0
\]

From this one finds the outgoing fluxes of \(E\) and \(L_z\) \(20\). This \(4\) also completely describes the radiation's interaction with the black hole, and thus encodes the rate at which the large black hole absorbs \(E\) and \(L_z\) \(21, 23\).

Teukolsky derived a "master equation" for black hole perturbations \(24\), and showed it separates by expanding in Fourier modes and spheroidal harmonics:

\[
4 = \frac{1}{(c^2 \sin^2) \int} \frac{\partial^Z X}{\partial t} R_{\text{in}}(z) S_{\text{in}} : (\hat{e}^m \hat{m}) : t
\]

(3)

(The parameter \(a = \gamma \pm M\) is the black hole spin per unit mass.) It is not necessary to expand in modes; one can leave the Teukolsky equation as a PDE coupled in \(t, r\), and (the dependence trivially decouples), and evolve initial data for \(4\). This works so well modeling source-free radiation \(27\) that it has been argued time domain methods may replace frequency decomposition for most applications \(28\). For point particle sources, frequency domain methods are currently considerably more accurate \(29\).

The radial functions \(R_{\text{in}}(z)\) are obtained by solving

\[
2 \frac{\partial^0}{\partial t^2} R_{\text{in}}(z) = V(z) R_{\text{in}}(z) = T_{\text{in}}(z) ; \quad (4)
\]

where prime denotes \(\partial = \partial_r = \partial_v^2 + 2M r^2 + \hat{e}\hat{v}, V(z)\) is a potential \(e.g., 20\), Eq. (4.3)) and \(T_{\text{in}}\) is a source constructed from the small body’s stress energy tensor. We build a Green’s function \(G_{\text{in}}(\beta z^2)\) from solutions to the homogeneous wave equation (setting \(T_{\text{in}} = 0\), using certain analytic transformations \(31\) that allow high numerical accuracy. We then find \(R_{\text{in}}(z)\) by integrating \(G_{\text{in}}(\beta z^2)\) over the source,

\[
R_{\text{in}}(z) = \int d\beta G_{\text{in}}(\beta z^2) T_{\text{in}}(z) \quad (5)
\]

Our particular interest is in the limits \(r! 1\) and \(r\) the event horizon; from the function in these limits, we extract the \(E\) and \(L_z\) fluxes mode-by-mode. Using many modes, we assemble the adiabatic rates of change \(\dot{E}_{Lz} \) and \(\dot{E}_{Lz} \).

For bound orbits, the source is nicely described by a harmonic expansion. The continuous frequency \(\nu\) goes over to a discrete set \(\nu_{n kn} = m + k + n r\), where \(m, k\) describe motion in \(l, i\) and \(r\) \(32,33\). Our modes become 4 index objects, \(\langle m kn \rangle\). There is no mode-mode coupling, so this problem is ideally suited to parallel computation: different modes can be sent to different processors with little communication cost. Figure \(\text{II}\) illustrates parallelization. The top panel shows the average CPU time per mode \(\langle m kn \rangle\) as a function of number of processors \(N\). We find nearly linear scaling in \(1/N\), demonstrating that we have little parallelization overhead. Typical CPU time per mode is quite short, so that \(\nu^{10} \nu^{10}\) modes can be computed in a reasonable time.

The bottom panel shows the convergence of energy flux from a strong field orbit. We write the orbit’s radial motion as \(r = p^m (1 + e \cos i)\); for this plot, we put eccentricity \(e = 0.5\) and semi-latus rectum \(p = 4\). The inclination \(i = 45\); the black hole’s spin is \(a = 0.5 M\). (Precise definitions of \(p, e, i\) and \(\nu\) can be found in \(33\); their key property for this paper is that they are simply related to the constants \(E, L_z, Q\).) We plot

\[
E_\nu = \sum X^l X^X X^X E_{m kn}
\]
where $E_{lm kn}$ is the modal energy flux. The cutoff values $N$ and $K$ (which formally are 1), are set large enough that neglected terms contribute a fractional amount less than $10^{-4}$ to the sum; details of this truncation will be presented elsewhere [34]. We generically find that $E_{lm kn}$ falls exponentially with $l$. The orbital parameter’s influence on the rate of this falloff, and on other convergence criteria, will be presented in [34]. Most importantly, the convergence of these fluxes is fairly quick for $e < 0.7$, a very astrophysically important range [5].

Mino’s adiabatic self force and $Q$’s evolution. Recent work by Mino [13] makes it possible to compute the time average $\Omega_{Q}$ of $Q$, allowing us to compute $\frac{\partial f}{\partial t}$ in Eq. (1) for generic orbits. The regularized self force $f^{\Omega}$ involves an integral over the orbiting body’s past worldline [11,12]. Mino shows that this integral can be replaced in the adiabatic limit by a relatively simple expression involving the difference between “retarded” and “advanced” forces: $f^{\Omega} = (f^{ret} - f^{adv})/2$. The “retarded” force depends on events on the orbiting body’s past lightcone; the “advanced” force depends on the future lightcone. Their difference removes the divergent point-particle self interaction; the remaining force drives the inspiral. This result is strikingly similar to Dirac’s result [10], and indeed reproduces the rule posited (without proof) by Gal’tsov [34].

The Carter constant is given by $Q = Q_{ab} p^{a} p^{b}$, where $Q_{ab}$ is a Killing tensor and $p^{a}$ is the small body’s 4-momentum. Taking a time derivative yields $\dot{Q} = 2Q_{ab} p^{a} p^{b} = 2Q_{ab} p^{a} p^{b}$, where $Q_{ab}$ is a Killing tensor and $p^{a}$ is the small body’s 4-momentum. We express the self force in terms of the radiative Green’s function using Mino’s result, and use the expansion of the radiative Green’s function in terms of modes [35]. Time averaging yields an expression of the form [16]

$$\dot{Q}_{Q_{lm kn}}(r) = X_{l} \times \int \int \int W_{l m kn}(r'; R_{lm kn}(r'));$$

(7)

The quantities $R_{lm kn}(r)$ come from $R_{lm kn}(r)$ by writing $R_{lm kn}(r) = l_{lm kn} R_{lm kn}(r)$, where $l_{lm kn}$ is the Wronskian which is independent of $r$ near the horizon and near infinity. The quantities $Q_{lm kn}(r)$ are computed via an integral similar to that in Eq. (6), but with the source $T_{lm kn}$ replaced by a new source built from the Killing tensor $Q_{ab}$ and the orbiting body’s 4-velocity, and evaluated at $t = t_{lm kn}$. Finally $W$ is the Wronskian which is independent of $r$ near the horizon and near infinity as both $R_{lm kn}(r)$ and $Q_{lm kn}(r)$ satisfy the homogeneous Teukolsky equation at those locations. Details of this calculation will be presented elsewhere [16]. Using this result it will be as straightforward to compute $\Omega_{Q}$ as it is currently to compute $\Omega_{H}$ and $\Omega_{H_{-1}}$.

Applications and future directions. Building inspiral waveforms requires us to compute backreaction effects upon a dense “grid” of orbits in parameter space. Each orbit is represented by a coordinate $(\varphi, \Omega, \rho)$; backreaction gives the tangent $(\delta \varphi, \delta \Omega, \delta \rho)$ to the inspiral trajectory. We flow along this tangent vector, building the trajectory $(\varphi(t); \Omega(t); \rho(t))$. Choosing initial conditions, we can likewise build the inspiral worldline $X^{a}(t) = [\varphi(t); \Omega(t); \rho(t)]$. To date, only circular $(e = 0)$ [36] and equatorial $(e = 0; \varphi = 180)$ [37] inspirals have been computed; the general case is under development [34], though we can present waveform “snapshots” (Fig. 2).

![Graph](image-url)
orbits, transferring the hole’s spin to the orbit, just as planetary tides can transfer angular momentum to a satellite. If the hole rotates rapidly and the orbit has shallow inclination, the tide can significantly prolong the inspiral \[36\].

*Strong field precessions:* The time to oscillate through \(\pi\) does not equal the time to move through 2 radians of \(\phi\). The mismatch between these timescales gives perihelion precession, a classic GR test. With black holes, this effect can amount to thousands of radians per orbit — the small body “whirls” many times near the hole before “zooming” out to large radius. This “zoom-whirl” behavior leaves a distinctive stamp on the waveform \[37\], seen in the short, high-frequency segments in Fig. 2.

*Spin-orbit coupling:* The black hole’s spin makes the space-time geometry oblate, and drags spacetime into co-rotation with it. This splits the \(c\tau\) corrections in the post-3.5-Newtonian phase of the Fourier transform of the waveform that correspond to : (t) give a phase error between 0.003 Hz and 0.03 Hz (the last year of inspiral \[18\]) of 0.2 cycles (after optimizing time-of-arrival and overall phase). For detection purposes, one needs phase coherence for only 3 weeks \[5\]; the phase error over this timescale will be even smaller.

Future work will develop ARR and waveforms with a focus upon templates for future GW searches. One goal is to use spectral methods for solving many of the formalism’s equations. Such methods are typically exponentially convergent in the number of basis functions. Since ARR requires many multipoles, each must be as accurate as possible. We are very encouraged by the work of Fujita and Tagoshi \[38\], who demonstrate that a particular set of basis functions allows modal solutions with essentially double precision accuracy.

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