Scattering Amplitudes and the Conservative Hamiltonian for Binary Systems at Third Post-Minkowskian Order

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We present the amplitude for classical scattering of gravitationally interacting massive scalars at third post-Minkowskian order. Our approach harnesses powerful tools from the modern amplitudes program such as generalized unitarity and the double-copy construction, which relates gravity integrands to simpler gauge-theory expressions. Adapting methods for integration and matching from effective field theory, we extract the conservative Hamiltonian for compact spinless binaries at third post-Minkowskian order. The resulting Hamiltonian is in complete agreement with corresponding terms in state-of-the-art expressions at fourth post-Newtonian order as well as the probe limit at all orders in velocity. We also derive the scattering angle at third post-Minkowskian order and find agreement with known results.

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Introduction.—The recent discovery of gravitational waves at LIGO/Virgo [1] has launched an extraordinary new era in astronomy, astrophysics and cosmology. Given expected improvements in detector sensitivity, high-precision theoretical predictions from general relativity will be crucial. Existing theory benchmarks come from a variety of approaches (see also Ref. [2] and references therein), including the effective one-body formalism [3], numerical relativity [4], the self-force formalism [5], and perturbative analysis using post-Newtonian (PN) [6–10], post-Minkowskian (PM) [11–13], and effective field theory (EFT) [14] methods.

The past decade has also witnessed immense progress in the study of scattering amplitudes, where understanding mathematical structures within gauge theory and gravity has yielded new physical insights and efficient methods for calculation. In particular, the Bern-Carrasco-Johansson (BCJ) color-kinematics duality and associated double copy construction [15] allow multiloop gravitational amplitudes to be constructed from sums of products of gauge-theory quantities. This has yielded a variety of new results in supergravity (see Ref. [16] for recent results). The BCJ construction is intimately tied to the Kawai-Lewellen-Tye (KLT) relations [17], which relate tree amplitudes of closed and open strings.

In this Letter, we apply modern amplitude methods to derive the classical scattering amplitude for two massive spinless particles at $\mathcal{O}(G^4)$ and to all orders in the velocity, i.e., at the third post-Minkowskian (3PM) order. We use generalized unitarity [18] to construct the corresponding two-loop integrand from tree amplitudes of gravitons and massive scalars, obtained straightforwardly from the double-copy construction. While the double copy introduces dilaton and antisymmetric tensor degrees of freedom [19], which are absent in pure Einstein gravity, we remove these unwanted states efficiently by restricting the state sums in unitarity cuts to gravitons alone. As we will show, we can calculate in strictly $D = 4$ dimensions for the classical dynamics, where spinor helicity variables [20,21] dramatically simplify the required tree amplitudes. The viability of working in $D = 4$ offers optimism for extending our results to higher orders.

Afterwards, we integrate the two-loop integrand via a procedure adapted from EFT, in which energy integrals are evaluated in the potential region via residues before performing spatial integrations [22]. Using EFT matching [22,23] we then derive the 3PM conservative Hamiltonian for compact spinless binaries. We show that the 4PN terms in our Hamiltonian are, up to a coordinate transformation, physically equivalent to corresponding terms in state-of-the-art results. We also verify that our result agrees in the probe limit with the Hamiltonian for a test body orbiting a Schwarzschild black hole to 3PM order. Finally, we derive a compact expression for the 3PM scattering angle in terms of amplitude data.

Double copy and unitarity.—Dynamics at 3PM order is encoded in the two-loop scattering amplitude for two
massive, gravitationally interacting scalars. Our calculation begins with a construction of the corresponding two-loop integrand via generalized unitarity. Because we are interested in classical scattering, we need not assemble the full quantum-mechanical integrand. Rather, as emphasized in Refs. [22–24], the classical potential only receives contributions with a single on-shell matter line per loop and with no gravitons starting and ending on the same matter line. For this reason we focus solely on the unitarity cuts shown in Fig. 1.

We obtain the tree amplitudes in the unitarity cuts via two methods. In the first approach, we work in general D space-time dimensions. Exploiting color-kinematics duality [15], we derive gravitational amplitudes straightforwardly from simpler gauge-theory amplitudes by replacing color factors with corresponding kinematic factors. For the unitarity cuts of the classical limit of the two-loop scattering amplitude, the reference momenta that complicate unitarity cuts of the classical limit of the two-loop scatter-

![Diagram](https://example.com/diagram.png)

**FIG. 1.** Unitarity cuts needed for the classical scattering amplitude. The shaded ovals represent tree amplitudes while the exposed lines depict on-shell states. The wiggly and straight lines denote gravitons and massive scalars, respectively.

...mass $m_1$ while legs $2^x, 3^y$ have mass $m_2$. All momenta in each tree amplitude are taken to be outgoing. The sum runs over graviton states for legs 5, 6, 7, and 8, where the minus signs on the labels indicate reversed momenta.

The four-point gravity tree amplitudes needed in the cuts are obtained from gauge-theory ones via the field-theory limit of KLT relations [17],

$$M_4(1, 2, 3, 4) = -i s_{12} A_4(1, 2, 3, 4) A_4(1, 2, 4, 3),$$

where the $A_4$ are tree-level color-ordered gauge-theory four-point amplitudes and $s_{ij} = (p_i + p_j)^2$, working in mostly minus metric signature throughout. Strictly speaking, the KLT relations apply only to massless states. However, they can be applied here by interpreting the scalar masses, in the sense of dimensional reduction, as extradimensional momentum components. While we have not included coupling constants, these are easily restored at the end of the calculation by including an overall factor of $(8\pi G)^3$, where $G$ is Newton’s constant.

In terms of the spinor-helicity conventions of Ref. [21], the independent tree-level gauge-theory amplitudes needed in Eq. (1) are

$$A_4(1^+, 2^+, 3^+, 4^+) = \frac{m_1^2 [23]}{(23) t_{12}},$$

$$A_4(1^+, 2^+, 3^-, 4^+) = \frac{\langle 3 \rangle [12]^2}{t_{23} t_{12}},$$

$$A_4(1^-, 2^-, 3^+, 4^+) = i \frac{\langle 12 \rangle^4}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle},$$

$$A_4(1^-, 2^+, 3^-, 4^+) = i \frac{\langle 13 \rangle^4}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle},$$

where $t_{ij} = 2p_ip_j$ and the ± denote gluon helicities.

The dilaton and antisymmetric tensor states are removed from unitarity cuts by correlating the gluon helicities on both sides of the double copy. The unwanted states correspond to one gluon in the double copy of positive helicity and the other of negative helicity. An internal graviton state is obtained by taking the corresponding gluons in the KLT formula in Eq. (2) to be of the same helicity.

Using spinor evaluation techniques, it is straightforward to obtain a compact expression for the iterated two-particle cut in Eq. (1) (e.g., see Ref. [26]). Imposing cuts on the matter lines, as indicated in the first unitarity cut of Fig. 1, further simplifies it and gives $C^{\text{H-cut}}$. We find

$$C^{\text{H-cut}} = 2i \left(\frac{1}{(p_5 - p_8)^2} + \frac{1}{(p_5 + p_7)^2}\right) \times \left(s_{23} m_4^2 + \frac{1}{s_{23}} \sum_{i=1}^{6} \left(\mathcal{E}_i^2 + O_1^2 + 6O_2^2 \mathcal{E}_1^2\right)\right).$$

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where we have defined
\[
E_1^2 = \frac{1}{4} s_{23}^2 (t_{18} t_{25} - t_{12} t_{58})^2, \quad O_1^2 = E_1^2 - m_1^2 m_2^2 s_{23} t_{58}, \\
E_2^2 = \frac{1}{4} s_{23}^2 [t_{17} t_{25} - t_{12} t_{57} - s_{23} (t_{17} + t_{57})]^2, \quad O_2^2 = E_2^2 - m_1^2 m_2^2 s_{23}^2 t_{57}.
\]

The simplicity of this expression is a reflection of the double-copy structure: the same building blocks appear in the simpler corresponding gauge-theory cut.

The spurious double pole in \( s_{23} \) can be explicitly cancelled by adding terms proportional to the Gram determinant formed from the five independent momenta at two loops that vanishes in \( D = 4 \). In fact, the expression derived from the \( D \)-dimensional approach is automatically free of such spurious singularities. While these Gram determinants contribute quantum mechanically, we have checked explicitly that they vanish in the classical limit. This is not accidental—such terms are of the wrong form to generate the required \( \log(s_{23}) \) needed to contribute to the classical 3PM amplitude (see Ref. [25] for details).

The remaining two independent generalized unitarity cuts in Fig. 2 are more complicated because they require five-point tree amplitudes with two massive scalar legs. The four-dimensional input gauge-theory amplitudes are simple to compute using modern methods (e.g., see Ref. [27]). For our \( D \)-construction we obtain a BCJ representation, allowing us to express the gravity cuts directly in terms of local diagrams. The particular representation was chosen such that we can ignore the reference momenta when projecting the internal states into gravitons. Further details will be given elsewhere [25].

To facilitate integration, we merge the cuts into a single integrand whose cuts match those in Fig. 1. This is achieved using an ansatz in terms of eight independent diagrams with only cubic vertices displayed in Fig. 2. The diagrammatic numerators are polynomials of the appropriate dimension exhibiting the symmetries of the corresponding diagram. Their coefficients are then fixed via the method of maximal cuts [28], whereby cuts of the integrand are constrained to match the known ones. This approach is sufficient for the two-loop problem.

Integration.—Our method of integration follows Ref. [22]. For convenience, we give a short summary here, leaving details to Ref. [25]. Terms in the integrand take the form,
\[
\mathcal{I} = \frac{\text{numerator graviton propagators}}{\prod_i \frac{1}{\omega_i^2 - k_i^2 - m_i^2}},
\]
where \( i \) labels each matter line, which has energy \( \omega_i \), spatial momentum \( k_i \), and mass \( m_i \). The matter propagators can be factored into particle and antiparticle poles, \( \omega_i \pm \sqrt{k_i^2 + m_i^2} \).

We then express the integrand as \( \mathcal{I} = \mathcal{N} \prod_i (1/z_i) \), i.e., in terms of the particle poles \( z_i = \omega_i - \sqrt{k_i^2 + m_i^2} \) and an effective numerator \( \mathcal{N} \) that absorbs the rest of the integrand.

Following the procedure outlined in Ref. [22], we first evaluate the energy integrals. At two loops, i.e., 3PM order, we integrate over two independent combinations of energies, \( \omega \) and \( \omega' \), in the potential region. As we will prove in detail in Ref. [25], the result is
\[
\tilde{\mathcal{I}} = \int \frac{d\omega d\omega'}{2\pi} \mathcal{I}(\omega, \omega') = \sum_{(i,j)} S_{ij} \text{Res} \mathcal{I}(\omega_i, \omega_j),
\]
where the sum runs over distinct pairings \((i, j)\) of matter poles and \( z_i = z_j = 0 \) when \( (\omega, \omega') = (\omega_{ij}, \omega_{ij}') \). Here \( S_{ij} \) is a calculable symmetry factor whose sign and magnitude depend on the topology of the cut graph. Note that the residue for an \((i, j)\) pairing will vanish if there are no values of \( \omega \) and \( \omega' \) for which \( z_i = z_j = 0 \).

The resulting quantity \( \tilde{\mathcal{I}} \) depends on two independent spatial loop momenta. To integrate over them we employ dimensional regularization to deal with ultraviolet divergences stemming from the renormalization of delta function contact interactions, which do not contribute classically. Due to the localization on energy residues, \( \tilde{\mathcal{I}} \) is a complicated, nonpolynomial function of three-dimensional invariants involving square roots. Nevertheless, we can series expand \( \tilde{\mathcal{I}} \) in large \( m_{1,2} \), yielding polynomials of kinematic invariants, which we can integrate at each order. After expanding, nearly all the spatial integrals are simple bubbles for which there are known analytic expressions [29]. The remaining integrals are evaluated via integration-by-parts identities [30].

For diagrams free from infrared (IR) singularities generated by iterations of lower-loop graviton exchanges, we have checked that our integrated results accord with several standard methods in the Feynman integral literature, including the Mellin-Barnes representation [29,31], numerical integration via sector decomposition [32], and differential equations [33] derived through integration-by-parts reduction [30,34]. The system of differential equations omits integrals lacking support on the matter pole residues that produce the classical contributions.

Amplitude and potential.—The integration procedure outlined above yields the conservative, i.e., real component...
of the 3PM amplitude generated by potential gravitons order by order in the large-mass expansion. Combining an explicit evaluation of this amplitude up to 7PN order with knowledge of the polar structure of individual integrals and exact, manifestly relativistic analytic results for certain graph topologies, we conjecture a full, all orders in velocity expression for the conservative 3PM amplitude (whose uniqueness will be discussed in Ref. [25]):

\[
M_3 = \frac{\pi G^2 \nu \gamma m^4 \log q^2}{6 \gamma^2 \xi} \left[ 3 - 6 \nu + 206 \nu \sigma - 54 \sigma^2 + 108 \nu \sigma^2 + 48 \nu (3 + 12 \sigma^2 - 4 \sigma^4) \arcsinh \sqrt{\frac{\sigma^2}{\gamma^2} - 1} \right. \\
- \frac{18 \nu \gamma (1 - 2 \sigma^2)(1 - 5 \sigma^2)}{(1 + \gamma)(1 + \sigma)} + \frac{8 \pi^2 G^3 \nu^4 \gamma^6}{\gamma^4 \xi} \left[ 3 \nu (1 - 2 \sigma^2)(1 - 5 \sigma^2) F_1 - 32 m^2 \nu (1 - 2 \sigma^2)^3 F_2 \right],
\]

where the log scale dependence is absorbed into a delta-function ultraviolet counterterm. Here we use center-of-mass coordinates where the incoming and outgoing particle momenta are \( \pm p \) and \( \pm (p - q) \), respectively. We emphasize that \( M_3 \) includes the nonrelativistic normalization factor, 1/4E1E2, where \( E_{1,2} = \sqrt{p^2 + m_{1,2}^2} \). We also define the total mass \( m = m_1 + m_2 \), the symmetric mass ratio \( \nu = m_1 m_2 / m^2 \), the total energy \( E = E_1 + E_2 \), the symmetric energy ratio \( \xi = E_1 E_2 / E^2 \), the energy-mass ratio \( \gamma = E / m \), and the relativistic kinematic invariant \( \sigma = p_1 \cdot p_2 / m_1 m_2 \). We emphasize that Eq. (8) is not valid for \( m_{1,2} \to 0 \) since quantum terms of order \( |q| / m_{1,2} \) are dropped, as will be elaborated on in Ref. [25]. Also, note that the arcsinh factor is proportional to the sum of particle rapidities, \( \text{arctanh} |p| / E_{1,2} \).

Equation (8) only includes \( q \)-dependent terms that persist in the classical limit. The \( \log q^2 \) term ultimately feeds into the conservative Hamiltonian through the Fourier transform \( \log q^2 |_{\text{FT}} = -1/2 \pi |r|^3 \). The IR-divergent contributions, parametrized by \( F_1 = \int k_1 / X_1^3 Y_1 X_2 \) and \( F_2 = \int k_1 k_2 / X_1^3 Y_1 X_2^3 Y_2 X_3 \) in the notation described in Eq. (12) of Ref. [22], will cancel in the EFT matching.

The Hamiltonian is extracted from the amplitude via EFT methods developed in Refs. [22,23,35] (see Ref. [13] for another approach). Consider massive spinless particles interacting via the center-of-mass Hamiltonian

\[
H(p, r) = \sqrt{p^2 + m_1^2} + \sqrt{p^2 + m_2^2} + V(p, r),
\]

\[
V(p, r) = \sum_{i=1}^{\infty} c_i (p^2) \left( \frac{G}{|r|} \right)^i,
\]

where \( r \) is the distance vector between particles and \( i \) labels PM orders. The above Hamiltonian is in a gauge in which terms involving \( p \cdot r \) or time derivatives of \( p \) are absent. We then compute the scattering amplitude of massive scalars, \( M^{(\text{EFT})} = \sum_{i=1}^{\infty} M_i^{(\text{EFT})} \), where \( M_3^{(\text{EFT})} \) comes from diagrams with two or fewer loops that depend on \( c_1 \), \( c_2 \), and \( c_3 \). In Ref. [22], the coefficients \( c_1 \) and \( c_2 \) were extracted analytically to all orders in velocity. Inserting these into \( M_3^{(\text{EFT})} \) effectively implements the subtraction of iterated contributions. By equating \( M_3^{(\text{EFT})} = M_3 \), we solve for the 3PM coefficient \( c_3 \).

The main result of the present Letter is the 3PM potential, encapsulated in the coefficients

\[
c_1 = \frac{\nu^2 m^2}{\gamma^2 \xi} (1 - 2 \sigma^2), \quad c_2 = \frac{\nu^2 m^2}{\gamma^2 \xi} \left[ \frac{3}{4} (1 - 5 \sigma^2) \right. \\
\left. - \frac{4 \nu \sigma (1 - 2 \sigma^2)}{(1 + \gamma)(1 + \sigma)} - \frac{\nu^2 (1 - \xi)(1 - 2 \sigma^2)^2}{2 \gamma \xi^2 \sigma^2}, \right] \\
- \frac{4 \nu (3 + 12 \nu^2 - 4 \sigma^4) \arcsinh \sqrt{\frac{\sigma^2}{\gamma^2} - 1}}{\sqrt{\gamma^2 - 1}} + \frac{3 \nu (1 - 2 \sigma^2)(1 - 5 \sigma^2)}{2 (1 + \gamma)(1 + \sigma)} - \frac{3 \nu \sigma (7 - 20 \sigma^2)}{2 \gamma \xi^2} + \frac{\nu^2 (3 + 8 \nu - 3 \xi - 15 \sigma^2 - 80 \nu \sigma^2 + 15 \xi \sigma^2)(1 - 2 \sigma^2)}{4 \gamma \xi^3} + \frac{2 \nu^2 (3 - 4 \xi \sigma)(1 - 2 \sigma^2)^2}{2 \gamma \xi^2} + \frac{\nu^4 (1 - 2 \xi)(1 - 2 \sigma^2)^3}{2 \gamma \xi^4}. \tag{10}
\]

where, for convenience, the expressions for \( c_1 \) and \( c_2 \) in Ref. [22] are reproduced here with slightly different normalization and in our current notation. As emphasized in Ref. [22], the cancellation of IR divergences between \( M_3^{(\text{EFT})} \) and \( M_3 \) depends critically on \( c_1 \) and \( c_2 \) and thus provides a nontrivial check of our calculation.

Consistency checks.—Our results pass several nontrivial albeit overlapping consistency checks (see Ref. [25] for details). First and foremost, we have verified that the 4PN terms in our Hamiltonian are equivalent to known results up to a canonical coordinate transformation,

\[
(r, p) \to (R, P) = (A r + B p, C p + D r) \tag{11}
\]

\[
A = 1 - \frac{G \nu}{2 |r|} + \ldots, \quad B = G (1 - 2 / \nu) \frac{p \cdot r}{4 m |r|} + \ldots,
\]

\[
C = 1 + \frac{G \nu}{2 |r|} + \ldots, \quad D = -\frac{G \nu}{2 |r|^3 |p \cdot r| + \ldots.
\]
treatments, see Ref. [36,37]). To derive this coordinate transformation we generate an ansatz for $A$, $B$, $C$, and $D$ and constrain it to preserve the Poisson brackets, i.e., $\{ r, p \} = \{ R, P \} = 1$ with all other brackets vanishing, in the spirit of Ref. [38]. We verify that within this space of canonical transformations exists a subspace that maps our Hamiltonian in Eq. (10) to the one in the literature, e.g., as canonical transformations exists a subspace that maps our Hamiltonian, we obtain the 3PM-accurate classical treatment, see Ref. [36,37]). To derive this coordinate transformation we generate an ansatz for $A$, $B$, $C$, and $D$ and constrain it to preserve the Poisson brackets, i.e., $\{ r, p \} = \{ R, P \} = 1$ with all other brackets vanishing, in the spirit of Ref. [38]. We verify that within this space of canonical transformations exists a subspace that maps our Hamiltonian in Eq. (10) to the one in the literature, e.g., as summarized in Eq. (8.41) of Ref. [10], up to the intersection of 3PM and 4PN accuracy.

Second, applying the methods of Ref. [22], we have checked that the full-theory amplitude $M_3$ in Eq. (8) is identical to the amplitude $M_3^{\text{EFT}}$ computed from the conservative Hamiltonian in Ref. [10] up to 4PN accuracy.

Third, we have extracted from our Hamiltonian the coordinate invariant energy of a circular orbit as a function of the period. Working at 2PN order—the highest order subsumed by 3PM which is relevant to a virialized system—we agree with known results [8].

Fourth, by solving the equations of motion derived from our Hamiltonian, we obtain the 3PM-accurate classical scattering angle in the center-of-mass frame and neglecting radiation,

$$2\pi\gamma = \frac{d_1}{J} + \frac{d_2}{J^2} + \frac{1}{J^3}( -4d_3 + \frac{d_4}{\pi} - \frac{d_5^3}{48\pi^3} ),$$

where $J = \sqrt{bp}$ is the angular momentum, $b$ is the impact parameter, and we have defined $d_1 = \frac{m\gamma_2}{\gamma q^2} M_1^1/p$, $d_2 = \frac{m\gamma_2}{\gamma q^2} M_2^1$, and $d_3 = \frac{m\gamma_2}{\gamma q^2} M_3^1 / \log q^2$, where the $q$ dependence cancels. The primed quantities denote the IR-finite parts of the nonrelativistically normalized amplitudes that enter the Hamiltonian coefficients as defined here and in Ref. [22], so

$$M_1^1 = -\frac{4\pi G\nu^2 m_2}{\nu^2 q^2} (1 - 2\sigma^2),$$

$$M_2^1 = -\frac{3\pi^2 G^2 \nu^2 m_1 m_2}{2\nu^2 \gamma q} (1 - 5\sigma^2),$$

and $M_3^1$ is the log $q^2$ term in Eq. (8). Truncated to 4PN order, Eq. (12) agrees with known results [39].

Last but not least, in the probe limit $m_1 \ll m_2$, our result exactly coincides with the Hamiltonian for a point particle in a Schwarzschild background to $O(G^0)$ and all orders in velocity, e.g., as given in Eq. (8) of Ref. [40].

Conclusions.—We have presented the 3PM amplitude for classical scattering of gravitationally interacting massive spinless particles. From this amplitude we have extracted the corresponding conservative Hamiltonian for binary dynamics to 3PM order.

The 3PM Hamiltonian in Eqs. (9) and (10) will be employed in a forthcoming paper [41] to compute approximants for the binding energy of binary systems moving on circular orbits and assess their accuracy against numerical-relativity predictions. This is relevant for understanding the usefulness of PM calculations when building accurate waveform models for LIGO/Virgo data analysis.

Our Letter leaves many avenues for future work, e.g., including obtaining higher orders in the PM expansion, incorporating spin [42], radiation [43], and finite-size effects, as well as connecting to other recent amplitude approaches [19,44] and the effective one-body formulation [3,12,13,45].

The simplicity of the 3PM amplitude in Eq. (8) and potential in Eq. (10) bodes well for future progress. Moreover, since the amplitude and EFT methods employed in this Letter are far from exhausted, we believe that the results we have reported mark only the beginning.

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