Iterative approach to arbitrary nonlinear optical response functions of graphene

F. Hipolito,1, ∗ Darko Dimitrovski,1 and T. G. Pedersen1,2, 1

1Department of Physics and Nanotechnology, Aalborg University, DK-9220 Aalborg Øst, Denmark
2Center for Nanostructured Graphene (CNG), DK-9220 Aalborg Øst, Denmark

Two-dimensional materials constitute an exciting platform for nonlinear optics with large nonlinearities that are tunable by gating. Hence, gate-tunable harmonic generation and intensity-dependent refraction have been observed in e.g. graphene and transition-metal dichalcogenides, whose electronic structures are accurately modelled by the (massive) Dirac equation. We exploit on the simplicity of this model and demonstrate here that arbitrary nonlinear response functions follow from a simple iterative approach. The power of this approach is illustrated by analytical expressions for harmonic generation and intensity-dependent refraction, both computed up to ninth order in the pump field. Moreover, the results allow for arbitrary band gaps and gating potentials. As illustrative applications, we consider (i) gate-dependence of third- and fifth-harmonic generation in gapped and gapless graphene, (ii) intensity-dependent refractive index of graphene up to ninth order, and (iii) intensity-dependence of high-harmonic generation.

Nonlinear optical (NLO) response encompasses a large class of light matter interactions [1–6], including processes such as harmonic generation and self-focusing of light, that have proven useful in a number of applications in nonlinear spectroscopy and in optoelectronic devices. Recent progress in the fabrication of 2D materials has produced a new fertile class of materials with large nonlinear susceptibilities [7]. Recent reports include measurements of high harmonic generation (HHG) [8, 9] and intensity dependent refractive index [10, 11] in graphene and in transition metal dichalcogenides (TMDs) [12]. In addition, it has been shown that the NLO response can be tuned by electrostatic doping [13–15] and significant progress has been made in measuring even-order NLO response in TMDs [16, 17]. Furthermore, the nonlinearities in 2D materials can be significantly enhanced by several mechanisms such as plasmons [18–20], polaritonic effects [21], and metasurfaces [22].

Compared with the linear response, calculations of NLO processes in crystals are significantly more complex. Whereas the linear response results from purely inter- or intraband processes, the NLO response contains not only these processes, but also mixed ones involving both inter- and intraband motion of electrons [23–26]. To circumvent this complexity, the NLO response has been characterized using several theoretical methods, each with its own merits and shortcomings: (i) perturbative expansion of the reduced density matrix [26–28]; (ii) time-dependent techniques [29–32]; (iii) Wannier representation [33]. The perturbative method offers a feasible approach to specific processes at a fixed frequency and power of the external field. The standard approach expands all matrix elements in unperturbed eigenstates, leading to increasingly complicated sum-over-states expressions for high-order processes. Still, within the perturbative regime, highly accurate results are obtained and in simple few-band systems such as the Dirac Hamiltonian, closed form solutions can often be found. These allow for characterization with respect to external parameters, for instance doping and temperature [26–28]. Yet, the growth in complexity associated with mixed inter- and intraband motion makes calculations extremely cumbersome beyond third order. The complicated nature of the general third-order response [26] testifies to this complexity. Methods (ii) and (iii) can be applied to study the NLO response at field strengths beyond the perturbative regime, as these intrinsically include contributions from all powers of the external field. But, in contrast to perturbative approaches, these methods rely extensively on numerical techniques for the integration of the equation of motion and for the Fourier transforms required to analyze the response in frequency domain, thus making the characterization of the NLO response with respect to external parameters an elaborate numerical process [29–32].

In the present letter, we study the (massive and massless) Dirac Hamiltonian as a model of graphene and TMDs. For this important class of materials, we bridge a key shortcoming found in perturbative techniques by evaluating the current density response via an iterative solution. This approach allows for the evaluation of arbitrarily high order response functions. As an illustration, we compute all response functions up to ninth order for (gapped) graphene [34]. The massive Dirac Hamiltonian [35] with a external vector potential \( \mathbf{A} = A_0 \sin(\omega t) \) reads, using the minimal coupling (velocity gauge), [36]

\[
H = v_F (\hbar \mathbf{\kappa} + e \mathbf{A}) \cdot \mathbf{\sigma} + \hbar \Delta \sigma_z ,
\]

where \( v_F \sim 10^6 \text{ m/s} \) is the Fermi velocity, \( \mathbf{\kappa} \) is the wavevector, \( \mathbf{\sigma} \) are the Pauli matrices, \( \Delta \geq 0 \) is the mass term, and \( e > 0 \). This model leads to a gapped band structure with energy gap \( E_g = 2\hbar \Delta \), and doping is including via a non-vanishing Fermi level \( \hbar \mu \). Hence, for pristine graphene, we take \( \Delta = 0 \). The key features of the electronic structure are shown in Fig. 1(a). The Dirac Hamiltonian has proven extremely useful for systems with threefold rotation symmetry. It allows for accurate analytic
characterization of several physical properties in graphene [37] and in the vast class of TMDs [38].

The time evolution of the wave function \( \psi \) governed by the unperturbed Hamiltonian, i.e., taking \( A = 0 \). We write the general wave function as

\[
\psi = [a(t)u_1 + b(t)u_2] \exp[i(\mathbf{k} \cdot \mathbf{r} + \epsilon t)],
\]

with energy dispersion \( \epsilon = \sqrt{\Delta^2 + k^2} \) (in frequency units), and \( k = v_F|\mathbf{k}| \). Furthermore, we focus on the response to a normally incident, linearly polarized monochromatic field, \( A = A_0 \sin(\omega t)/\hbar \) and \( \mathbf{A}(t) = v_F A_0 \sin(\omega t) \), which is related to the electric field \( \mathcal{E}(t) = \mathcal{E}_0 \cos(\omega t) \) by \( \mathcal{E}_0 = -\omega A_0 \). The time evolution of the coefficients follows from

\[
\dot{a}(t) = i\mathcal{A}(t)[a(t)F - b(t)G],
\]

and

\[
\dot{b}(t) = \mathcal{A}(t)G^\ast a(t) + [2\epsilon + \mathcal{A}(t)F]b(t),
\]

where \( F = \sqrt{\Delta^2 - \epsilon^2}/(\epsilon + G) \) and \( G = (\Delta \cos(\theta) - i\epsilon \sin(\theta)) / \epsilon \) arise from matrix elements of the velocity operator. The coefficients then determine the reduced density matrix, whose matrix elements read \( \rho_{11} = |a|^2 \), \( \rho_{22} = |b|^2 \) and \( \rho_{21} = a^\ast b \). In turn, their time evolution is governed by

\[
i\dot{\mathcal{N}} = -2\mathcal{A}(t)(P\mathcal{G} - \mathcal{G}^\ast P^\ast),
\]

\[
i\dot{\mathcal{P}} = -\mathcal{A}(t)\mathcal{G}^\ast \mathcal{N} + 2[\mathcal{A}(t)F + \epsilon]\mathcal{P},
\]

where \( \mathcal{N} = \rho_{22} - \rho_{11} \) and \( \mathcal{P} = \rho_{21} \) define the population difference and coherence, respectively. Finally, the current density is evaluated via the expectation value of the current density operator \( j = -e\mathbf{g}_x g_y (2\pi)^{-2} v_F \int \mathcal{J} d\mathbf{k} \), where the dimensionless integrand for the current density response is defined by \( \mathcal{J} = \sum_{mn} v_{nm} p_{nm} / v_F = \mathcal{N} + \mathcal{P}^\ast + \mathcal{P} + \mathcal{G}^\ast \mathcal{P}^\ast \) using the matrix elements \( v_{nm} \) of the velocity operator \( \mathbf{v}_x = v_F \mathbf{e}_x \). Here, \( g_x = 2 \) and \( g_y = 2 \) are spin and valley degeneracies, respectively.

The iterative sequence is found by considering the first- and second-order time derivatives of \( \mathcal{J} \) that read

\[
\mathcal{A}(t) \dot{\mathcal{J}} = -\epsilon \mathcal{N},
\]

\[
\ddot{\mathcal{J}} = 4\epsilon [\mathcal{A}(t) + \epsilon F] \mathcal{N} - 4\epsilon [\mathcal{A}(t)F + \epsilon] \mathcal{J}.
\]

Using a time-harmonic expansion for the integrand \( \mathcal{J} = \sum_n \mathcal{J}_n e^{-i\omega t} \) and for the population \( \mathcal{N} = \sum_n \mathcal{N}_n e^{-i\omega t} \), the dynamical equations can be cast as

\[
\mathcal{A}_0 [n - 1] \mathcal{J}_{n-1} - (n + 1) \mathcal{J}_{n+1} = 2i\epsilon n \mathcal{N}_n,
\]

\[
(4\epsilon^2 - n^2 \omega^2) \mathcal{J}_n = 4\epsilon^2 F \mathcal{N}_n - 2i\epsilon A_0 F (\mathcal{J}_{n-1} - \mathcal{J}_{n+1}) + 2i\epsilon A_0 (\mathcal{N}_{n-1} - \mathcal{N}_{n+1}),
\]

where \( n \in \mathbb{Z} \) defines the Fourier order. The final iterative series for the integrand is identified by making use of an expansion with respect to powers of the external field \( \mathcal{J}_n = \sum_{j \geq 0} \mathcal{J}_n^{(j)} \mathcal{A}_j^0 \) and collecting equal powers [39]

\[
\mathcal{J}_n^{(j)} = \Theta_{j,n+4} K_4 \mathcal{J}_n^{(j-2)} + \Theta_{j,n+2} [2 \Theta_{j,n} K_0 \mathcal{J}_n^{(j-2)} + \delta_{n,0} K_1 \mathcal{J}_n^{(j-1)}]
\]

where

\[
\Theta_{j,n+4} K_4 \mathcal{J}_n^{(j-2)} + \Theta_{j,n+2} [2 \Theta_{j,n} K_0 \mathcal{J}_n^{(j-2)} + \delta_{n,0} K_1 \mathcal{J}_n^{(j-1)}]
\]

\[
\mathcal{J}_n^{(j)} = \Theta_{j,n+4} K_4 \mathcal{J}_n^{(j-2)} + \Theta_{j,n+2} [\Theta_{j,n} K_0 \mathcal{J}_n^{(j-2)} + \delta_{n,0} K_1 \mathcal{J}_n^{(j-1)}]
\]

\[
\mathcal{J}_n^{(n)} = K_{-2} \mathcal{J}_{n-2} + K_{-1} \mathcal{J}_{n-1},
\]

which lends itself to a diagrammatic representation as illustrated in Fig. 1(b-c).

To apply the iterative solution for \( \mathcal{J}_n^{(j)} \) in practice, two seeds \( \mathcal{J}_0^{(0)} \) and \( \mathcal{J}_1^{(1)} \) are required that can easily be determined from low-order terms in Eq. (5b). Collecting all terms independent of the external field and making use of the equilibrium charge distribution \( \mathcal{N}_0 = f_1 - f_2 \) (the difference between Fermi functions), the first seed reads \( \mathcal{J}_0^{(0)} = F \mathcal{N}_0 \). The second seed involves the collection of linear terms in the external field and reads \( \mathcal{J}_1^{(1)} = 2\epsilon [G^2 \mathcal{N}_0 A_0 / (4\epsilon^2 - \omega^2)] \). All remaining terms of \( \mathcal{J}_n^{(j)} \) can be computed sequentially by evaluating all possible Fourier components \( n = \{0, 1, \ldots, j\} \), in increasing order, for any given response order \( j \) using Eq. (6). The solutions for all nonzero integrands up to fifth order are listed in this order in Tab. 1. The final response functions are obtained by integrating the desired \( \mathcal{J}_n^{(j)} \) over \( \mathbf{k} \)-space and the respective conductivities \( \sigma_n^{(j)}(\omega) \) then follow by writing

\[
\mathcal{J}(t) = \sum_{n,j} \sigma_n^{(j)}(\omega) \mathcal{E}_0 e^{-i\omega t} + \text{c.c.} / 2.
\]
TABLE I. Analytic expressions for current density integrands up to fifth order, where $J^{(j)}_n \equiv F^{(j)}_n |G|^2 N_0 A_0^j$.

| $(n, j)$ | $F^{(j)}_n$ |
|----------|-------------|
| (0, 0)   | $2e/[(4\epsilon^2 - \omega^2)]$ |
| (1, 1)   | $-2F/[4\epsilon^2 - \omega^2]$ |
| (2, 2)   | $6e^2 F/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (1, 3)   | $4i\epsilon[(13F^2 - 1)e^2 - (4F^2 - 1\omega^2)]/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (3, 3)   | $-4i\epsilon [(5F^2 - 1)e^2 + \omega^2]/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (0, 4)   | $-2F[2(13F^2 - 7)e^2 - (8F^2 - 5\omega^2)]/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (2, 4)   | $16e^2 F[8(8F^2 - 3)e^4 - 2(53F^2 - 28)e^2 \omega^2 + (27F^2 - 17)\omega^4]/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (4, 4)   | $-10e^2 F[(7F^2 - 3)e^2 + 5\omega^2]/[(4\epsilon^2 - 16\omega^2)(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (1, 5)   | $48ie^2[61F^4 - 32F^2 + 1]e^0 - 2(503F^4 - 334F^2 + 11)e^2 \omega^2 + 5(100F^4 - 77F^2 + 4)e^4 \omega^4$ |
|          | $-3(24F^4 - 21F^2 + 2)\omega^6]/[4(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |
| (3, 5)   | $-16ie^2[4(295F^4 - 186F^2 + 11)e^8 - (5235F^4 - 4166F^2 + 291)e^6 \omega^2 + (6095F^4 - 6303F^2 + 558)e^4 \omega^4 - (1440F^4 - 2569F^2 + 419)e^2 \omega^8 + 36(3 - 8F^2)\omega^8]/[4(4\epsilon^2 - 16\omega^2)(4\epsilon^2 - \omega^2)^2(4\epsilon^2 - \omega^2)^2]$ |
| (5, 5)   | $12ie^2[(21F^4 - 14F^2 + 1)e^4 - 5(1 - 7F^2)e^2 \omega^2 + 4\omega^4]/[(4\epsilon^2 - 25\omega^2)(4\epsilon^2 - 16\omega^2)(4\epsilon^2 - \omega^2)(4\epsilon^2 - \omega^2)]$ |

In most cases, the integration is straightforward, but can lead to cumbersome expressions, particularly whenever the difference between the Fourier $n$ and response order $j$ is large. The angular part of the integral depends exclusively on powers of $F$ and $|G|$, therefore it can be shown that due to the presence of full rotation symmetry in the effective Hamiltonian all even-order response functions vanish upon angular integration. Nonetheless, even-order integrands remains necessary to determine higher order non-vanishing odd integrands.

Now, we turn our attention to the conductivities computed within the iterative framework. At low temperatures the population difference becomes a step function $N_0 = -\Theta(\epsilon - |\mu|)$ and the lower limit $\Omega$ of the radial part of the integral is determined by the larger of the Fermi level and mass term, i.e. $\Omega \equiv \max(|\mu|, \Delta)$ [41]. In our explicit examples, we compute all conductivities up to ninth order [34]. Among these, we examine in detail third and fifth harmonic generation as well as intensity-dependent refraction through the optical Kerr effect including high-order terms. As demonstrated in recent experiments [13–15], valuable information can be extracted by varying the Fermi level via electrostatic gating. Hence, we apply the present results to study the doping dependence of these NLO processes.

The third harmonic generation (THG) conductivity reads

$$
\sigma_3^{(3)}(\omega) = \frac{-3i\sigma_3}{1024\pi} \left( \frac{2v_F}{3a_0\omega} \right)^4 \left[ \frac{56\Delta^2}{\omega^2} - \frac{48\Delta^4}{\omega^4} \right] \ln \frac{2\Omega - 3\omega}{2\Omega + 3\omega} \\
- \left( \frac{128\Delta^2}{\omega^2} - \frac{192\Delta^4}{\omega^4} \right) \ln \frac{2\Omega - 2\omega}{2\Omega + 2\omega} + \left( 17 + \frac{88\Delta^2}{\omega^2} - \frac{240\Delta^4}{\omega^4} \right) \ln \frac{2\Omega - \omega}{2\Omega + \omega}, (9)
$$

where we define the scale of the 2D nonlinear conductivities systematically by $\sigma_{j+1} \equiv 2(3\epsilon a_0^2/(4\hbar v_F))^{j+1} \sigma_1$ with $\sigma_1 = e^2/(4\hbar)$ and the carbon-carbon distance $a_0 \equiv 1.42$ Å sets the natural length scale for graphene. Throughout the letter, we consider exclusively electron doping $\mu > 0$, but results for hole doping $\mu < 0$ simply follow by replacing $\mu \rightarrow -\mu$. Taking the limit $\Delta \rightarrow 0$, one can verify that our expression reduces to previous results derived with the gapless Dirac Hamiltonian using velocity and length gauges [26–28]. The expression for $\sigma_3^{(5)}(\omega)$ is representative of the HHG conductivities, $\sigma_n^{(n)}(\omega)$, which are always composed by $n$ logarithmic divergences, whose amplitude is set by a polynomial prefactor with even powers of $\Delta/\omega$ as shown in Eq. S1a (see supplemental material). Note that the divergences found in the expressions are regularized by introducing a small broadening parameter $\omega \rightarrow \omega + i\eta$ and, unless stated otherwise, we use $\hbar\eta = 1$ meV. The THG conductivities for both gapped and gapless graphene assuming photon energies in the low and medium range are shown in Fig. 2. Given the rather small gaps that can be reliably generated in graphene [42, 43] (we take $E_g \equiv 250$ meV as a reference figure for our calculations) and considering photon energies $h\omega > 150$ meV, our results show that the response of gapped and gapped systems are generally similar but deviate whenever $\mu \lesssim \Delta$.

The fifth harmonic conductivity in gapless graphene also lends itself to a closed form expression [34], with $n = 5$ logarithmic divergences

$$
\sigma_5^{(5)}(\omega) = \frac{9i\sigma_5}{65536\pi} \left( \frac{2v_F}{3a_0\omega} \right)^8 \left[ 4925 \ln \frac{2\mu - 5\omega}{2\mu + 5\omega} + \frac{17}{\omega} \right],
$$
FIG. 2. THG conductivity in gapped ($E_g = 250$ meV, black) and gapless (red) graphene at photon energies $\hbar\omega = \{150, 400, 800\}$ meV.

\[ -16384 \ln \frac{2\mu - 4\omega}{2\mu + 4\omega} + 19359 \ln \frac{2\mu - 3\omega}{2\mu + 3\omega} \]
\[ - 9216 \ln \frac{2\mu - 2\omega}{2\mu + 2\omega} + 1266 \ln \frac{2\mu - \omega}{2\mu + \omega}, \]

where the general expression valid for $\Delta \neq 0$ is given in Eq. (S3) [34]. The fifth-order response of graphene is highly sensitive to the ratio between photon energy and doping level. This is illustrated in the contour plot in Fig. 3 of the amplitude of the fifth harmonic conductivity as function of these parameters. It shows that this response function can be tuned over several orders of magnitude by tuning either parameter, while highlighting the five resonances present in the fifth harmonic response. Moreover, we find that the nonlinear conductivities of graphene are regular in the limit of vanishing doping $\mu \to 0$

\[ \lim_{\mu \to 0} \frac{\sigma_n^{(j)}}{\sigma_j} = q_{nj} \left( \frac{2v_F}{3a_0}\right)^{2(j-1)}, \]

where the coefficients $q_{nj}$ are rational numbers. The complete list for all coefficients is found in Tab. S1 in supplemental material. For third and fifth harmonic generation, the coefficients read $q_{33} = 3/512$ and $q_{55} = -45/65536$, respectively.

The iterative approach can also readily be used to evaluate conductivities beyond harmonic generation such as the optical Kerr conductivity of graphene [34]

\[ \sigma_1^{(3)}(\omega) = \frac{9i\sigma_0}{256\pi} \left( \frac{2v_F}{3a_0}\right)^4 \left\{ \frac{12\mu\omega}{4\mu^2 - \omega^2} \right. \]
\[ \left. - 11 \ln \frac{\omega - 2\mu}{\omega + 2\mu} + 16 \ln \frac{2\omega - 2\mu}{2\omega + 2\mu} \right\}, \]

This expression is representative of high-order contributions to any Fourier order $\sigma_n^{(j \geq n)}$ [34]. These expressions contain $(j + n)/2$ logarithmic divergences, rather than $j = n$ found in the $n^{th}$ harmonic conductivities, and also contain an additional rational function with $(j + n - 2)/2$ polynomial divergences that strongly enhance the nonlinear resonances, see Fig. S1 in supplemental material. In Fig. 4, we show the conductivities $\sigma_1^{(j)}$ contributing to the first harmonic current in doped graphene up to ninth order at $E_0 = 2V/\mu m$. Note that the field intensity considered in Fig. 4 matches the upper limit of the perturbative regime when considering THz radiation [9]. Inspection of Fig. 4 defines the regime, where the perturbative approach breaks down, namely the frequency range, in which terms $\sigma_1^{(j)}E_0^{j-1}$ cease to decrease as the order $j$ is increased. Hence, for the parameters in Fig. 4 the non-perturbative region can be estimated as $\hbar\omega \lesssim 60$ meV. Manifestations from higher than Kerr terms should be detectable as higher order terms introduce additional resonances that are highly sensitive to both the Fermi level and the magnitude of the external field.

In Fig. 5, we plot the relative amplitude of the Fourier components of the radiated intensity $I_n(\omega) = \mu_0 c |J_n(\omega)|^2/8$ with $J_n(\omega) = \sum_j \sigma_n^{(j)}(\omega)E_0^j$, with respect to the incident intensity in vacuum $I_0 = \epsilon_0 c E_0^2/2$ [44, 45] considering all contributions up to $j = 15$. Note that the analytic expressions for the conductivities are limited to ninth order, hence all data plotted in Fig. 5 were integrated.
FIG. 5. Intensity of HHG Fourier components in doped graphene ($\hbar\mu = 250$ meV) normalized by $I_0 = \varepsilon_0 c_0 E_0^2/2$. The incident photon energy is $\hbar\omega = 100$ meV and colors red to blue represent increasing field strengths.

numerically using $\hbar\eta = 10$ meV. Results shown in Figs. 2 to 5 demonstrate that the approach presented in this letter can be used to readily characterize harmonic response of graphene, including the effects from higher order terms, at arbitrary doping level and photon frequency, without requiring the complex numerical calculations found in time-dependent techniques.

In summary, we introduce and apply an iterative approach to the calculation of NLO response of systems described by the massive Dirac Hamiltonian. The iterative nature allows for analytical evaluation of high-order response functions, and we derive for the first time all nonlinear conductivities of (gapped) graphene up to ninth order. The NLO response of doped graphene reveals an intricate interplay between doping, photon energy and the intensity of the external electric field.

The authors acknowledge Alireza Taghizadeh for many helpful comments. This work was supported by the QUSCOPE center sponsored by the Villum Foundation, and TGP is supported by the CNG center under the Danish National Research Foundation, project DNRF103.

* fh@nano.aau.dk
† tgp@nano.aau.dk

[1] R. W. Boyd, *Nonlinear Optics*, 3rd ed. (Elsevier Science Publishing Co Inc, 2008).

[2] Y. R. Shen, *The Principles of Nonlinear Optics* (Wiley-Interscience, 2002).

[3] P. A. Franken and J. F. Ward, Rev. Mod. Phys. 35, 23 (1963).

[4] N. Bloembergen, Rev. Mod. Phys. 54, 685 (1982).

[5] V. M. Axt and S. Mukamel, Rev. Mod. Phys. 70, 145 (1998).

[6] M. G. Kuzyk, K. D. Singer, and G. I. Stegeman, Adv. Opt. Photon. 5, 4 (2013).

[7] A. Autere, H. Jussila, Y. Dai, Y. Wang, H. Lipsanen, and Z. Sun, Adv. Mater. 30, 1705963 (2018).

[8] N. Yoshikawa, T. Tamaya, and K. Tanaka, Sci. 356, 736 (2017).

[9] H. A. Hafez, S. Kovaliev, J. C. Deinert, Z. Mics, B. Green, N. Awari, M. Chen, S. Germanskii, U. Lehnert, J. Teichert, Z. Wang, K. J. Tielrooij, Z. Liu, Z. Chen, A. Narita, K. Müllen, M. Bonn, M. Gensch, and D. Turchinovich, Nature 561, 507 (2018).

[10] G.-K. Lim, Z.-L. Chen, J. Clark, R. G. S. Goh, W.-H. Ng, H.-W. Tan, R. H. Friend, P. K. H. Ho, and L.-L. Chua, Nat. Photon. 5, 554 (2011).

[11] M. Mohsin, D. Neumaier, D. Schall, M. Otto, C. Matheisen, A. Lena Giesecke, A. A. Sagade, and H. Kurz, Sci. Rep. 5, 10967 (2015).

[12] H. Liu, Y. Li, Y. S. You, S. Ghimire, T. F. Heinz, and D. A. Reis, Nat. Phys. 13, 262 (2016).

[13] G. Soavi, G. Wang, H. Rostami, D. G. Purdie, D. De Fazio, T. Ma, B. Luo, J. Wang, A. K. Ott, D. Yoon, S. A. Bourelle, J. E. Muench, I. Goykhman, S. Dal Conte, M. Celelbrano, A. Tomadin, M. Polini, G. Cerullo, and A. C. Ferrari. Nat. Nanotechnol. 13, 583 (2018).

[14] T. Jiang, D. Huang, J. Cheng, X. Fan, Z. Zhang, Y. Shan, Y. Yi, Y. Dai, L. Shi, K. Liu, C. Zeng, J. Z. Ji, J. E. Sipe, Y.-R. Shen, W.-T. Liu, and S. Wu, Nat. Photon. 12, 430 (2018).

[15] Y. Zhang, D. Huang, Y. Shan, T. Jiang, Z. Zhang, K. Liu, J. Cheng, J. E. Sipe, W.-T. Liu, and S. Wu, arXiv 1812.11306 (2018).

[16] G. Wang, X. Marie, I. Gerber, T. Amand, D. Ladgaré, L. Bouet, M. Vidal, A. Balocchi, and B. Urbaszek, Phys. Rev. Lett. 114, 097403 (2015).

[17] A. Siyutikoki, L. Karvonen, H. Rostami, A. Autere, S. Mehravar, A. Lombardo, R. A. Norwood, T. Hasan, N. Peyghambarian, H. Lipsanen, K. Kieu, A. C. Ferrari, M. Polini, and Z. Sun, Nat. Commun. 8, 893 (2017).

[18] J. D. Cox, A. Marin, and F. J. G. de Abajo, Nat. Commun. 8, 14380 (2017).

[19] D. Kundys, B. Van Duppen, O. P. Marshall, F. Rodriguez, I. Torre, A. Tomadin, M. Polini, and A. N. Grigorenko, Nano Lett. 18, 282 (2018).

[20] Z. Wang, Z. Dong, H. Zhu, L. Jin, M.-H. Chiu, L.-J. Li, Q.-H. Xu, G. Eda, S. A. Maier, A. T. S. Wee, C.-W. Qiu, and J. K. W. Yang, ACS Nano 12, 1859 (2018).

[21] D. S. Wild, E. Shahmoon, S. F. Yelin, and M. D. Lukin, Phys. Rev. Lett. 121, 123606 (2018).

[22] G. Rosolen, L. J. Wong, N. Rivera, B. Maes, M. Soljačić, and I. Kaminer, Light: Sci. Appl. 7, 64 (2018).

[23] C. Aversa and J. E. Sipe, Phys. Rev. B 52, 14636 (1995).

[24] T. G. Pedersen, Phys. Rev. B 92, 235432 (2015).

[25] A. Taghizadeh, F. Hipolito, and T. G. Pedersen, Phys. Rev. B 96, 195413 (2017).

[26] H. P. Hipolito, A. Taghizadeh, and T. G. Pedersen, Phys. Rev. B 98, 205420 (2018).

[27] J. L. Cheng, N. Vermeulen, and J. E. Sipe, New J. Phys. 16, 053014 (2014).

[28] H. Rostami and M. Polini, Phys. Rev. B 93, 161411 (2016).

[29] T. Tamaya, A. Ishikawa, T. Ogawa, and K. Tanaka, Phys. Rev. Lett. 116, 016601 (2016).

[30] L. A. Chizhova, F. Libisch, and J. Burgdörfer, Phys. Rev. B 95, 085436 (2017).

[31] D. Dimitrovski, T. G. Pedersen, and L. B. Madsen, Phys. Rev. A 95, 063420 (2017).

[32] S. A. Mikhailov, Phys. Rev. B 95, 085432 (2017).

[33] F. Catoire, H. Bachau, Z. Wang, C. Blaga, P. Agostini, and
L. F. DiMauro, Phys. Rev. Lett. **121**, 143902 (2018).

[34] See Supplemental Material at [URL will be inserted by publisher] for the full expressions for gapped graphene and for additional information.

[35] G. W. Semenoff, Phys. Rev. Lett. **53**, 2449 (1984).

[36] N. M. R. Peres, Rev. Mod. Phys. **82**, 2673 (2010).

[37] A. H. Castro Neto, F. Guinea, N. M. R. Peres, K. S. Novoselov, and A. K. Geim, Rev. Mod. Phys. **81**, 109 (2009).

[38] D. Xiao, G.-B. Liu, W. Feng, X. Xu, and W. Yao, Phys. Rev. Lett. **108**, 196802 (2012).

[39] Note: it is sufficient to consider \( n \geq 0 \), as the terms for \( J^{(j)}_n \) can be immediately obtained from \( J^{(j)}_n \) by means of the replacement \( \omega \rightarrow -\omega \).

[40] Note: the discrete unit step function is defined as \( \Theta_{i,j} = 1 \) for \( i \geq j \), \( \Theta_{i,j} = 0 \) for \( i < j \).

[41] Note: the integral over the wavevector has been replaced by an integration over energy.

[42] S. Y. Zhou, G.-H. Gweon, a. V. Fedorov, P. N. First, W. A. de Heer, D.-H. Lee, F. Guinea, A. H. Castro Neto, and A. Lanzara, Nat. Mater. **6**, 916 (2007).

[43] C. R. Woods, L. Britnell, A. Eckmann, R. S. Ma, J. C. Lu, H. M. Guo, X. Lin, G. L. Yu, Y. Cao, R. V. Gorbachev, A. V. Kretinin, J. Park, L. A. Ponomarenko, M. I. Katsnelson, Y. N. Gornostyrev, K. Watanabe, T. Taniguchi, C. Casiraghi, H.-j. J. Gao, A. K. Geim, and K. S. Novoselov, Nat. Phys. **10**, 451 (2014).

[44] T. Stauber, N. M. R. Peres, and A. K. Geim, Phys. Rev. B **78**, 085432 (2008).

[45] F. Hipolito and V. M. Pereira, 2D Mater. **4**, 021027 (2017).
We provide supplemental material for: i) the role of high order contributions to the Kerr conductivity; ii) the impact of doping on different Fourier components of the radiated field; iii) expressions for all nonlinear conductivities in (gapped) graphene up to ninth order.

**INTERPLAY BETWEEN LOGARITHMIC AND POLYNOMIAL PARTS**

The polynomial resonances found in higher order contributions to a given Fourier component, i.e. $j > n$, play an important role in the overall response of the system, by enhancing significantly the $(j + n - 2)/2$ lowest resonances. To illustrate this, we plot in Fig. S1 the contributions that arise from the logarithmic divergences and from the polynomial resonances in the optical Kerr conductivity $\sigma_n^{(3)}(\omega)$. The effect is two-fold, first and most striking is the dramatic increase in the amplitude of the lowest resonance by virtue of the polynomial resonance. Second is the destructive combination of both contributions for small doping $\mu < \omega/2$.

**HIGHER ORDER TERMS AND ROLE OF DOPING**

Doping plays an important role in the optical response of graphene. At linear order its effect is well known, primarily blocking the real part of the response for $\omega < 2\mu$ and sets the amplitude of the Drude peak. Yet, as discussed above and on the main paper (see Fig. 2) its role in the nonlinear response can be drastic. In Fig. S2, we plot the amplitude of the Fourier components of the radiated field at two different doping levels $h\mu = \{250, 312.5\}$ meV. Despite the relatively small change in doping, the nonlinear response changes drastically, particularly the Fourier components $n = \{3, 5, 7\}$, stressing once more the importance of the interplay between doping level and photon energy.

**CONDUCTIVITIES UP TO NINTH ORDER FOR GAPPED SYSTEMS**

The general expressions for the $n^{th}$ harmonic and higher order corrections $j > n$ in gapped graphene can be expressed in the form

$$\sigma_n^{(j)}(\omega) = \frac{i\sigma_n}{\pi} \left( \frac{2v_F}{3\epsilon_0\omega} \right)^2 \left( \frac{n+1}{2} \right)^{j-1} \sum_{m=1}^{(n+j)/2} \ln \left( \frac{2\Omega - m\omega}{2\Omega + m\omega} \right)$$

(S1a)
\[
\times \sum_{k=0}^{(j+1)/2} c_{km}^{(nj)} \left( \frac{\Delta}{\omega} \right)^{2k} + \frac{P_{nj}}{Q_{nj}} \right) (j > n),
\]
(S1b)

where \( c_{km}^{(nj)} \) are constant coefficients, \( P_{nj} \) and \( Q_{nj} \) are polynomial functions in \( \omega \), and recall the systematic definition of the nonlinear conductivity scales introduced in the main text \( \sigma_{j>1} \equiv 2\beta e a q / (\hbar v_F) \). As shown in Eq. (11) in the main text, all nonlinear conductivities are regular at vanishing doping. The complete list of non-zero coefficients \( q_{nj} \), up to ninth order, can be found in Tab. S1.

We provide the full expressions for the optical Kerr effect, fifth, seventh and ninth harmonic conductivities, namely Eqs. (S2), (S3), (S4) and (S5). For \( j > n \), we provide the coefficients \( c_{km}^{(nj)} \) in Tab. S2 and the respective polynomial functions \( P_{nj} \) and \( Q_{nj} \), found in Eqs. (S6) and (S7). Results for gapless graphene can be expressed directly from those of gapped graphene, by taking the limit \( \Delta \to 0 \) and replacing \( \Omega \) by \( \mu \).

The optical Kerr conductivity in gapped graphene reads
\[
\sigma_1^{(3)}(\omega) = \frac{9i\sigma_0}{256\pi} \left[ \frac{2v_F}{3q_0 a} \right]^4 \left[ \frac{4\lambda \Delta^2 \omega^4 + 8\lambda^2 \Gamma \Delta^2 \omega^6 + 48\lambda^4 (9\Omega^2 - 2\omega^2)}{\omega^6 \mu (4\mu^2 - \omega^2)} \right] + 16 + \frac{32\Delta^2 - 48\Delta^4}{\omega^2} \ln \frac{2\omega - 2\Omega}{2\omega + 2\Omega} - \frac{11 + 56\Delta^2}{\omega^2} - \frac{52835\Delta^4}{\omega^4} \ln \frac{2\omega - 2\Omega}{2\omega + 2\Omega}.
\]
(S2)

The fifth, seventh and ninth harmonic conductivities read
\[
\sigma_5^{(5)}(\omega) = \frac{9i\sigma_0}{65536\pi} \left[ \frac{2v_F}{3q_0 a} \right]^8 \left[ \frac{4925 + 3788\Delta^2}{\omega^2} \omega^6 - \frac{2640\Delta^4}{\omega^4} + \frac{320\Delta^6}{\omega^6} \right] \ln \frac{2\omega - 5\Omega}{2\Omega + 5\Omega} - \frac{16384}{2\Omega + 4\Omega} \frac{2\omega - 4\Omega}{2\omega + 4\Omega} \ln \frac{2\omega - 5\Omega}{2\Omega + 5\Omega} - \frac{9216}{2\Omega + 2\Omega} \frac{2\omega - 2\Omega}{2\omega + 2\Omega} \ln \frac{2\omega - 5\Omega}{2\Omega + 5\Omega} - 911408\Delta^2 \omega^4 - \frac{268040\Delta^4}{\omega^6} + \frac{107520\Delta^8}{\omega^8} \ln \frac{2\omega - 6\Omega}{2\Omega + 6\Omega} - \frac{584481}{2\Omega + 2\Omega} \frac{2\omega - 6\Omega}{2\omega + 6\Omega} \ln \frac{2\omega - 6\Omega}{2\Omega + 6\Omega} - \frac{386421}{2\Omega + 3\Omega} \frac{2\omega - 6\Omega}{2\omega + 3\Omega} \ln \frac{2\omega - 6\Omega}{2\Omega + 3\Omega} - \frac{104061}{2\Omega + 3\Omega} \frac{2\omega - 2\Omega}{2\omega + 3\Omega} \ln \frac{2\omega - 2\Omega}{2\Omega + 3\Omega}.
\]
(S3)

\[
\sigma_7^{(7)}(\omega) = \frac{9i\sigma_0}{2936012800\pi} \left[ \frac{2v_F}{3q_0 a} \right]^{12} \left[ \frac{5126492 + 211408\Delta^2}{\omega^2} \omega^6 - \frac{1680672\Delta^4}{\omega^4} + \frac{224000\Delta^6}{\omega^6} - \frac{860\Delta^8}{\omega^8} \right] \ln \frac{2\omega - 7\Omega}{2\Omega + 7\Omega} + \frac{2712268}{2\Omega + 4\Omega} \frac{2\omega - 4\Omega}{2\omega + 4\Omega} \ln \frac{2\omega - 7\Omega}{2\Omega + 7\Omega} - \frac{584481}{2\Omega + 4\Omega} \frac{2\omega - 4\Omega}{2\omega + 4\Omega} \ln \frac{2\omega - 7\Omega}{2\Omega + 7\Omega} - \frac{386421}{2\Omega + 3\Omega} \frac{2\omega - 4\Omega}{2\omega + 3\Omega} \ln \frac{2\omega - 7\Omega}{2\Omega + 3\Omega} - \frac{1127808}{2\Omega + 3\Omega} \frac{2\omega - 2\Omega}{2\omega + 3\Omega} \ln \frac{2\omega - 7\Omega}{2\Omega + 3\Omega}.
\]
(S4)

\[
\sigma_9^{(9)}(\omega) = \frac{27i\sigma_0}{10522669875200\pi} \left[ \frac{2v_F}{3q_0 a} \right]^{16} \left[ \frac{811948939 + 369488004\Delta^2}{\omega^2} \omega^6 - \frac{18232040\Delta^4}{\omega^4} + \frac{32327360\Delta^6}{\omega^6} - \frac{1191680\Delta^8}{\omega^8} + \frac{21504\Delta^{10}}{\omega^{10}} \right] \ln \frac{2\omega - 9\Omega}{2\Omega + 9\Omega} + \frac{5905580032}{2\Omega + 8\Omega} \frac{2\omega - 8\Omega}{2\omega + 8\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 8\Omega} + \frac{18559737203}{2\Omega + 7\Omega} \frac{2\omega - 7\Omega}{2\omega + 7\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 7\Omega} + \frac{32794509312}{2\Omega + 6\Omega} \frac{2\omega - 6\Omega}{2\omega + 6\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 6\Omega} + \frac{35426172500}{2\Omega + 5\Omega} \frac{2\omega - 5\Omega}{2\omega + 5\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 5\Omega} + \frac{23729834424}{2\Omega + 4\Omega} \frac{2\omega - 4\Omega}{2\omega + 4\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 4\Omega} + \frac{9439968636}{2\Omega + 3\Omega} \frac{2\omega - 3\Omega}{2\omega + 3\Omega} \ln \frac{2\omega - 9\Omega}{2\Omega + 3\Omega}.
\]
(S5)
TABLE S1. Coefficients $q_{nj}$ for nonlinear conductivities of graphene at the zero doping limit.

| $n$ | $j = 1$ | $j = 2$ | $j = 3$ | $j = 4$ | $j = 5$ | $j = 6$ | $j = 7$ | $j = 8$ | $j = 9$ |
|-----|---------|---------|---------|---------|---------|---------|---------|---------|---------|
| 1   | 45/256  | 1629/16384 | 1708371/10485760 | 41182970949/134217728000 |
| 3   | 3/512   | -351/32768 | -30475581/83886800 | -1174923089433/134217728000 |
| 5   | -45/65536 | 777122073/262144000 | -1205821487213379/164416176800000 |
| 7   |         | 7371/10485760 | -1251787974784857/46036680704000 |
| 9   |         |         | -3971997/751619276800 |

The polynomial functions of the family $P_{nj}$ read

$$P_{15} = -27 \left[ -320\Delta^6 \left( 6860\mu^8 - 9795\mu^6\omega^2 + 3149\mu^4\omega^4 - 216\mu^2\omega^6 - 16\omega^8 \right) + 48\Delta^4 \left( 4340\mu^6\omega^2 - 6469\mu^4\omega^4 + 2297\mu^2\omega^6 - 240\mu^8 \right) 
- 4\Delta^2 \left( 4\mu^8 + 433\mu^6\omega^2 - 59\mu^4\omega^4 \right) - 127\mu^6\omega^2 + 119\mu^8 + 5\mu^8 \omega^2 \right] \ln \frac{2\Omega - 2\omega}{2\Omega - 2\Omega + \omega}.$$

(S6a)

$$P_{17} = 81 \left[ 8960\Delta^6 \left( 220725120\mu^{16} - 108571760\mu^{14}\omega^2 + 1817194584\mu^{12}\omega^4 - 1439294370\mu^{10}\omega^6 + 506082893\mu^8\omega^8 - 77041635\mu^6\omega^{10} 
+ 3399312\mu^4\omega^{12} + 20160\mu^2\omega^{14} + 20736\omega^{16} \right) - 1280\Delta^6 \mu^2 \omega^2 \left( 24602816\mu^{14} - 121233232\mu^{12}\omega^2 + 209561345\mu^{10}\omega^4 - 161682431\mu^8\omega^6 
+ 507404056\mu^6\omega^8 - 862529394\mu^4\omega^{10} + 300768\mu^2\omega^{12} + 27648\omega^{14} \right) + 96\Delta^4 \mu^4 \omega^4 \left( 8962240\mu^{12} - 4433104\mu^{10}\omega^2 + 77619556\mu^8\omega^4 
- 60976999\mu^6\omega^6 + 22423706\mu^4\omega^{10} - 3822343\mu^2\omega^{12} + 24422\omega^{14} \right) - 16\Delta^2 \mu^6 \omega^6 \left( 3214720\mu^{10} - 1125392\mu^8\omega^2 + 964020\mu^6\omega^4 
- 250267\mu^4\omega^6 + 27383\mu^2\omega^{10} + 2700\omega^{12} \right) + \mu^8 \omega^8 \left( - 4269376\mu^{10} + 16069552\mu^8\omega^2 - 16274748\mu^6\omega^4 + 7469521\mu^4\omega^6 
- 1748206\mu^2\omega^{10} + 135657\omega^{12} \right) \right] \ln \frac{2\Omega - \omega}{2\Omega + \omega}.$$

(S6b)

$$P_{19} = 27 \left[ 4608\Delta^10 \left( 187256945075200\mu^{26} - 2325179092531546\mu^{24}\omega^2 + 11687240067038592\mu^{22}\omega^4 - 3122675397633056\mu^{20}\omega^6 
+ 48860030668451540\mu^{18}\omega^8 - 46478022674034201\mu^{16}\omega^{10} + 26989978924387373\mu^{14}\omega^{12} - 9364370277141650\mu^{12}\omega^{14} 
+ 1846375279461405\mu^{10}\omega^{16} - 181530911486796\mu^8\omega^{18} + 5430656668160\mu^6\omega^{20} + 169113692160\mu^4\omega^{22} - 1582266112\mu^2\omega^{24} 
+ 3224862720\omega^{26} \right) - 8690\Delta^8 \mu^2 \omega^2 \left( 20712412399104\mu^{24} - 2571207237387284\mu^{22} + 1292922374791414\mu^{20} 
- 3455335155552654\mu^{18} - 54074328703809484\mu^{16} - 51499663244720339\mu^{14} + 29884780216243249\mu^{12} - 1037651228740229\mu^{10} 
+ 204600002411288\mu^8\omega^6 - 20081754607804\mu^6\omega^8 + 677007583506\mu^4\omega^{10} + 10226995200\mu^2\omega^{12} + 35473489921\mu^2\omega^{24} + 3200\Delta^6 \mu^2 \omega^4 \left( 31160248073984\mu^2 - 39236536545360\mu^2\omega^2 
+ 197440066745984\mu^4\omega^2 - 527854548357440\mu^6\omega^4 + 826471420779144\mu^8\omega^6 - 7868266992735685\mu^{10}\omega^8 
+ 4573719204746129\mu^{12}\omega^{10} - 158728528597115\mu^{14}\omega^{12} + 311808889389989\mu^{16}\omega^{14} - 2978876192190843\mu^{18}\omega^{16} + 1718919063654594\mu^{20}\omega^{20} 
- 598137631786373\mu^{22}\omega^{22} + 12129145525643\mu^{24}\omega^{24} - 1313199585948\mu^{24} + 58525286400\mu^{24} \right) \right] \ln \frac{2\Omega - \omega}{2\Omega + \omega}.$$

(S6c)

$$P_{23} = 9 \left[ 1600\Delta^6 \left( 140\mu^6 - 495\mu^4\omega^2 + 439\mu^2\omega^4 - 72\omega^6 \right) + 48\Delta^4 \left( 100\mu^6 + 387\mu^4\omega^2 - 353\mu^2\omega^4 + 576\omega^6 \right) 
+ 4\Delta^2 \left( 430\mu^6 - 100\mu^4\omega^2 + 387\mu^2\omega^4 \right) + \mu^2 \omega^2 \left( 2572\mu^4 - 7231\mu^2\omega^2 + 819\omega^4 \right) \right] \ln \frac{2\Omega - \omega}{2\Omega + \omega}.$$

(S6d)
\[ P_{39} = 9 \left[ 64512A^{10} \left( 390753515495424^{12} \mu_{16} - 764147029967564^{14} \mu_{18}^{2} + 8041258112024658^{24} \mu_{4}^{4} - 475750643401113088^{8} \mu_{8}^{22} \right) + 175458495956524208^{10} \omega_{16}^{8} - 4223937579523353929^{8} \mu_{8}^{10} + 67519250001676414493^{8} \mu_{8}^{10} + 7145124722011657943^{14} \mu_{14}^{14} + 14985036538666293^{12} \mu_{12}^{12} - 208666654322865033^{10} \mu_{10}^{10} + 50589562622389564^{10} \mu_{10}^{20} - 6804986918332432^{6} \mu_{6}^{22} + 232023022041600^{24} \omega_{16}^{24} + 423767116800^{26} \mu_{26} + 1209323520000^{28} - 8960 \Delta^{8} \mu_{8}^{12} + 7907441 \omega_{12}, \right] \]

\[ P_{7} = -81 \left[ -6270A^{\Delta} \left( 1232^{10} \mu_{10}^{10} - 17016^{8} \mu_{12}^{8} + 7937^{6} \mu_{14}^{6} + 15997^{4} \mu_{6}^{4} + 10592^{2} \mu_{8}^{2} - 16200^{10} \right) + 44800 \Delta^{6} \omega_{6} \left( 3184^{10} - 4088^{8} \mu_{8}^{2} + 20755^{6} \mu_{10}^{6} - 39600^{4} \mu_{12}^{4} + 28395^{2} \mu_{8}^{2} - 43290^{10} \right) + 96 \Delta^{4} \left( -16340^{8} + 2304744^{6} \mu_{6}^{6} - 1120179^{4} \mu_{8}^{4} + 2277951^{2} \mu_{10}^{2} - 18591684^{2} \mu_{8}^{8} + 2764800^{10} + 16 \Delta^{2} \mu_{12}^{2} - 103184^{8} - 1342344^{6} \mu_{6}^{6} + 557810^{4} \mu_{8}^{4} - 790744^{2} \mu_{10}^{2} - 60390^{8} \right) + \mu_{2}^{2} \left( 1148188^{8} - 1880209^{6} \mu_{8}^{6} + 8225540^{4} \mu_{8}^{4} - 13934838^{2} \mu_{6}^{6} + 21260700^{10} \right) \right], \]

\[ P_{59} = 81 \left[ 193536^{10} \left( 95335871488^{24} - 3474667149056^{22} \omega_{22} + 53617091508608^{8} \mu_{8}^{4} + 45932503949968^{6} \mu_{6}^{6} + 240612655101334^{10} \mu_{10}^{10} + 8003718707211601^{14} + 17005840241750655^{12} \mu_{12}^{12} - 22640508661620861^{10} \mu_{10}^{10} + 1800553692912743^{8} \mu_{16}^{8} - 77893353191986^{6} \mu_{18}^{6} + 151814283623100^{4} \mu_{4}^{4} - 7805478432000^{2} \mu_{2}^{2} - 51438240000^{24} \right) + 62720^{\Delta} \left( -5934326323^{2} + 2188129607792^{22} \mu_{22}^{2} - 33670540569792^{20} \mu_{20}^{2} + 2885724233539312^{18} \mu_{18}^{6} + 1512519741686932^{16} \mu_{16}^{6} + 530567679412877^{14} \mu_{14}^{14} - 10709167127166907^{12} \mu_{12}^{12} + 1427723613182807^{10} \mu_{10}^{10} - 11375757423293642^{8} \mu_{8}^{8} + 4904774424609732^{6} \mu_{6}^{6} - 967438068232300^{4} \mu_{4}^{4} + 52236757440000^{2} \mu_{2}^{2} + 293932800000^{0} - 4480^{6} \omega_{6}^{4} - 2021600253184^{2} + 73794708322048^{22} \mu_{22}^{2} - 114112798916774^{20} \mu_{20}^{4} + 98024214792568^{18} + 5153815716433896^{16} \mu_{16}^{8} + 17223243800739083^{14} \mu_{14}^{14} + 3682413326318672^{12} \mu_{12}^{12} + 4947167054689431^{10} - 39766769343865879^{8} \mu_{8}^{8} + 17490666450972488^{6} \mu_{6}^{6} - 3512352132150000^{4} \mu_{4}^{4}. \right] \]
while the respective \(Q_{ij}\) polynomials read

\[Q_{15} = 512\omega^6 \mu^3 \left(4\mu^2 - 4\omega^2\right) \left(4\mu^2 - \omega^2\right), \quad (S7a)\]

\[Q_{17} = 16384\omega^7 \mu^3 \left(4\mu^2 - 9\omega^2\right) \left(4\mu^2 - 4\omega^2\right) \left(4\mu^2 - \omega^2\right)^2, \quad (S7b)\]

\[Q_{19} = 655360\omega^9 \mu^3 \left(4\mu^2 - 16\omega^2\right) \left(4\mu^2 - 9\omega^2\right)^2 \left(4\mu^2 - 4\omega^2\right)^3, \quad (S7c)\]

\[Q_{35} = 1024\omega^6 \mu \left(4\mu^2 - 9\omega^2\right) \left(4\mu^2 - 4\omega^2\right) \left(4\mu^2 - \omega^2\right)^2, \quad (S7d)\]

\[Q_{37} = 16384\omega^7 \mu ^3 \left(4\mu^2 - 16\omega^2\right) \left(4\mu^2 - 9\omega^2\right)^2 \left(4\mu^2 - 4\omega^2\right)^2 \left(4\mu^2 - \omega^2\right)^2, \quad (S7e)\]

\[Q_{39} = 104857600\omega^9 \mu^3 \left(4\mu^2 - 25\omega^2\right) \left(4\mu^2 - 16\omega^2\right)^2 \left(4\mu^2 - 9\omega^2\right)^3 \left(4\mu^2 - 4\omega^2\right)^3 \left(4\mu^2 - \omega^2\right)^3, \quad (S7f)\]

\[Q_{57} = 819200\omega^7 \mu \left(4\mu^2 - 25\omega^2\right) \left(4\mu^2 - 16\omega^2\right) \left(4\mu^2 - 9\omega^2\right) \left(4\mu^2 - 4\omega^2\right) \left(4\mu^2 - \omega^2\right), \quad (S7g)\]

\[Q_{59} = 1835008000\omega^9 \mu^3 \left(4\mu^2 - 36\omega^2\right)^2 \left(4\mu^2 - 25\omega^2\right)^2 \left(4\mu^2 - 16\omega^2\right)^2 \left(4\mu^2 - 9\omega^2\right)^2 \left(4\mu^2 - 4\omega^2\right)^2 \left(4\mu^2 - \omega^2\right)^2, \quad (S7h)\]

\[Q_{79} = 10276044800\omega^9 \mu \left(4\mu^2 - 49\omega^2\right)^2 \left(4\mu^2 - 36\omega^2\right)^2 \left(4\mu^2 - 25\omega^2\right)^2 \left(4\mu^2 - 16\omega^2\right)^2 \left(4\mu^2 - 9\omega^2\right)^2 \left(4\mu^2 - 4\omega^2\right)^2 \left(4\mu^2 - \omega^2\right). \quad (S7i)\]
| (n, m, j) | k = 0   | 1       | 2       | 3       | 4       | 5       |
|-----------|--------|---------|---------|---------|---------|---------|
| (1, 5, 1) | 10701  | 32768   | 13437   | 45327   | 10701   | 32768   |
| (1, 5, 2) | -45/32 | -153/32 | 1161/32 | -1395/32| -45/32  | -153/32 |
| (1, 5, 3) | 38637  | 32768   | 13437   | 45327   | 10701   | 32768   |
| (1, 5, 4) | 24894  | 32768   | 13437   | 45327   | 10701   | 32768   |
| (1, 5, 5) | 97972  | 32768   | 13437   | 45327   | 10701   | 32768   |

**TABLE S2. Coefficients \( c(nj) \) for nonlinear response functions.**