Wang, Yadong; Iyikanat, Fadil; Bai, Xueyin; Hu, Xuerong; Das, Susobhan; Dai, Yunyun; Zhang, Yi; Du, Luojun; Li, Shisheng; Lipsanen, Harri; García De Abajo, F. Javier; Sun, Zhipei

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Published in:
Nano Letters

DOI:
10.1021/acs.nanolett.2c02711

Published: 09/11/2022

Document Version
Publisher's PDF, also known as Version of record

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Please cite the original version:
Wang, Y., Iyikanat, F., Bai, X., Hu, X., Das, S., Dai, Y., Zhang, Y., Du, L., Li, S., Lipsanen, H., García De Abajo, F. J., & Sun, Z. (2022). Optical Control of High-Harmonic Generation at the Atomic Thickness. Nano Letters, 22(21), 8455-8462. https://doi.org/10.1021/acs.nanolett.2c02711

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Optical Control of High-Harmonic Generation at the Atomic Thickness

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Cite This: Nano Lett. 2022, 22, 8455−8462

ABSTRACT: High-harmonic generation (HHG), an extreme nonlinear optical phenomenon beyond the perturbation regime, is of great significance for various potential applications, such as high-energy ultrashort pulse generation with outstanding spatio-temporal coherence. However, efficient active control of HHG is still challenging due to the weak light–matter interaction displayed by currently known materials. Here, we demonstrate optically controlled HHG in monolayer semiconductors via the engineering of interband polarization. We find that HHG can be efficiently controlled in the excitonic spectral region with modulation depths up to 95% and ultrafast response speeds of several picoseconds. Quantitative time-domain theory of the nonlinear optical susceptibilities in monolayer semiconductors further corroborates these experimental observations. Our demonstration not only offers an in-depth understanding of HHG but also provides an effective approach toward active optical devices for strong-field physics and extreme nonlinear optics.

KEYWORDS: All-Optical Control, Static High-Harmonic Generation, Transient High-Harmonic Generation, Interband Carrier Transition, Electronic States, Real-Time Quantitative Theory of Nonlinear Optics

Recently, high-harmonic generation (HHG) in solid materials has attracted significant interest in strong-field and attosecond physics. While HHG based on atomic gases has formed the basis of attosecond science, the required complex setups with vacuum operation environments and delicate optics limit different practical applications. The emerging field of solid-state HHG naturally holds a series of advantageous features, such as a compact configuration to simplify the operation conditions, the requirement of lower pulse intensities (few μJ in solids vs few mJ in gases), and the involvement of strong electron interactions. The prospect to open possibilities extending beyond gas-phase HHG is thus strong motivation to further develop the field of extreme solid-state photonics. Since the pioneering observation of HHG in a bulk ZnO crystal, solid-state HHG has been demonstrated in a wide range of materials, including large-bandgap dielectrics, semiconductors, semimetals, topological insulators, and an ε-near-zero material. These results have revealed that HHG not only provides an effective method to probe the electronic structure in solids within the regime of strong interaction with light fields but also constitutes a promising platform for novel spectroscopy and microscopy techniques in advanced photonics and optoelectronics.

Active control over the extreme nonlinear processes involved in HHG would be a concrete step toward various applications. Currently, solid-state HHG is commonly understood as the result of two prominent physical mechanisms: interband polarization and light-driven intraband currents. Both mechanisms offer plausible opportunities to control HHG by manipulating the interband polarization or intraband current. For example, the HHG efficiency in bulk ZnO has been tuned by incoherently driving photoexcited carriers. Electrically controlled HHG in carbon nanotubes has also been reported by tuning the electron and hole doping levels. In addition, solid-state HHG has been engineered with symmetric or artificial structures, such as metallic plasmons and dielectric nanostructures. These abilities to tune or enhance HHG emphasize the potential of actively controlling HHG. The superior features of all-optical control (e.g., ease of operation, high speed, and low-power consumption) are expected to facilitate ultrafast optoelectronic applications, such as optical coding and data processing. Moreover, measurements of HHG in the time domain could play a fundamental role in understanding the physical mechanisms underlying the
extreme optics of solid-state materials.\textsuperscript{6,39} However, thus far, efficient control of HHG in solids remains elusive.

Recently, two-dimensional layered materials (2DLMs), which hold the advantages of strong optical nonlinear response and automatic fulfillment of the phase matching condition, have been found to serve as a powerful platform for HHG without the need to account for propagation effects, and thus, they are appealing for extreme nonlinear optics.\textsuperscript{12,14,40–45} Because of the strong interband oscillation strength in these materials, HHG in 2DLMs can be dramatically and resonantly enhanced by electronic states.\textsuperscript{1} All-optical active control of HHG from 2DLMs thus appears as an exciting opportunity.

In this work, we demonstrate an efficient optical control of HHG in monolayer MoS\textsubscript{2} and few-layer graphene. All-optical control is realized by tuning the interband polarization strength via the valence carrier density, achieving a high modulation depth of \(\sim 95\%\) in monolayer MoS\textsubscript{2} at an incident pulse energy of 250 \(\mu\)J and an ultrafast response time within the subpicosecond regime in few-graphene layer. Further, we present a real-time quantitative theory of nonlinear optical susceptibilities illustrating the underlying carrier dynamics. Our experimental and theoretical results show that interband polarization constitutes an effective method to modulate HHG in 2DLMs.

The monolayer MoS\textsubscript{2} flake used in our experiment is grown by chemical vapor deposition (CVD), and its quality is confirmed by photoluminescence and Raman spectroscopies (see Figure S1 in the Supporting Information). The nonlinear optical response in monolayer MoS\textsubscript{2} is analyzed at room temperature when illuminated with femtosecond seed pulses (\(\sim 150\) fs, 2 kHz, photon energy tunable in a range of \(\sim 0.52\) to 0.78 eV). To modulate the nonlinear optical response of the seed light in monolayer MoS\textsubscript{2}, we employ the control light with a photon energy \(\hbar\omega\) of \(\sim 3.1\) eV, as shown in Figure 1a. Figure 1b shows the HHG spectra obtained with and without the control light when the incident seed photon energy \(\hbar\omega\) is \(\sim 0.57\) eV with an average power of \(\sim 80\) \(\mu\)W and a delay time \(\Delta\tau \approx 0.4\) ps. In the HHG spectrum, the third-, fourth-, and fifth-harmonic generations (HGs) are observed at the photon energies of \(\sim 1.71, 2.29,\) and 2.85 eV, respectively. With the control light (incident power of \(\sim 100\) nW), the fourth- and fifth-harmonic signals significantly decrease, while the third-harmonic signal remains constant (see time-resolved 3\textsuperscript{rd}-HG in Figure S2). This indicates that the harmonic signal is not affected by the interband carrier density when the photon energy is below the optical bandgap of monolayer MoS\textsubscript{2}.\textsuperscript{34,35} To discuss the optical modulation of HHG in more detail, we focus on the fourth- and fifth-harmonic signals below.

First, we study the static fourth-order HHG (i.e., the control light is off). Figure 2a shows the fourth-harmonic spectrum when the seed photon energy \(\hbar\omega_{\text{seed}}\) is set to \(\sim 0.71\) eV (wavelength \(\lambda_{\text{seed}} = 1760\) nm) with an average power \((P_{\text{seed}})\) of 45 \(\mu\)W. The output photon energy is centered at \(\hbar\omega_{\text{4th}} \approx 2.84\) eV \((\lambda_{\text{4th}} \approx 437\) nm), confirming the fourth-harmonic generation process. Remarkably, the fourth-harmonic output power follows a power law \(P_{\text{4th}}\) with an index \(\alpha\) of 3 when the average seed light power is \(P_{\text{seed}} \leq 60\) \(\mu\)W (Figure 2b). This power dependence shows a lower slope than the power law of \(P_{\text{th}}\) that is typical of the perturbative regime (i.e., with a fourth-harmonic nonlinear polarization \(P_{\text{4th}}(\omega) \approx \varepsilon_{\text{opt}}(\omega)\chi^{(4)}(E(\omega))\), where \(\varepsilon_{\text{opt}}\) is the permittivity of free space, \(\chi^{(4)}\) corresponds to the fourth-order nonlinear optical susceptibility, and \(E(\omega)\) denotes the electric field of the incident seed-light beam, respectively). Figure 2c shows the peak intensities of the static 4\textsuperscript{th}-HG spectra via tuning the seed photon energy from 0.60 to 0.74 eV with \(P_{\text{seed}}\) fixed at 45 \(\mu\)W. The results show a power variation \((P_{\text{4th}})/P_{\text{max}}\) around 16 with a prominent peak at \(\sim 2.85\) eV and a small peak at \(\sim 2.43\) eV. We assign the peak at 2.85 eV to the C exciton formed by the band-nesting effect in monolayer MoS\textsubscript{2}.\textsuperscript{14,45} The peak at 2.43 eV could be caused by the other resonantly excited exciton states (such as the 3s exciton).\textsuperscript{46}

To achieve optically controlled 4\textsuperscript{th}-HG, we now introduce the control light with a photon energy \((\hbar\omega_{\text{ctrl}})\) of 3.1 eV and an average power \((P_{\text{ctrl}})\) of 100 nW. Figure 2d shows the transient fourth-harmonic signal \((\Delta P_{\text{4th}})\) generated from 2.7 to 2.87 eV versus the delay time between the control and seed pulses. We apply the definition \(\Delta P_{\text{4th}} = (P_{\text{4th}} - P_{\text{4th}})/P_{\text{max}}\) where \(P_{\text{4th}}\) and \(P_{\text{4th}}\) denote the fourth-harmonic signal power with and without the control light, respectively. Note that the controlled 4\textsuperscript{th}-HG at other wavelength regions is beyond our detection ability due to the relatively weak fourth-harmonic signal (Figure 2c). Nevertheless, the results demonstrate a broadband control of 4\textsuperscript{th}-HG around the band-nesting region of monolayer MoS\textsubscript{2}. Figure 2e shows \(\Delta P_{\text{4th}}\) at the 4\textsuperscript{th}-HG energy of 2.85 eV as a function of delay time \((\Delta\tau)\) between the control and seed pulses. We find that \(\Delta P_{\text{4th}}\) first decreases within 100 fs due to optical bleaching and subsequently starts to recover at \(\sim 30\) ps. The biexponential fitting of the recovering trace yields fast and slow time constants of \(\tau_1 \approx 2.2 \pm 0.16\) ps and \(\tau_2 \approx 57.6 \pm 10\) ps, respectively. We note that the carrier dynamics are independent of the control light power. Further, we examine the control power dependence of \(\Delta P_{\text{4th}}\) at the 4\textsuperscript{th} harmonic power \((\Delta\tau) \approx 0.4\) ps (Figure 2f), from which we find that...
ΔP_{4th} decreases linearly and sharply to a value of −65% with a control power of ∼150 nW and then slowly drops to −95% when the control power is increased to ∼500 nW (corresponding to a pulse energy of 250 pJ). This indicates full control over the 4th-HG intensity with the control light. The linear modulation dependence on the control light power can be attributed to the fact that the control light linearly excites the electrons in the valence band when its power is small.

Now, we focus on the 5th-HG in monolayer MoS₂. Figure 3 shows the static 5th-HG spectrum with the seed light power fixed at P₀ = 45 μW (I₀ ~ 0.65 TW/cm²) during experiments in (a) and (c)–(f).
seed photon energy ($\hbar \omega_0$) is 0.565 eV with an average power ($P_0$) of 50 μW. The peak output photon energy of the 5th-HG is at 2.80 eV, confirming its 5th-HG nature. When increasing the seed light power, the intensity of the 5th-HG signal increases nonlinearly (Figure 3b). It follows a power index of $\alpha$ $\approx$ 3.83 when the incident optical power is $\leq$ 80 μW, then starts being saturated. Similar to the 4th-HG, we find that the 5th-HG signal shows a large deviation in power dependence from the $P_0^\alpha$ behavior expected in the perturbative regime (i.e., with a fifth-harmonic nonlinear polarization $p^{(5)}(\omega) \approx \epsilon_0 \chi^{(5)}(\omega_0)\vec{E}(\omega_0)$, where the $\chi^{(5)}$ corresponds to the fifth-order nonlinear optical susceptibility). To investigate such deviation, we measure the wavelength-dependence of the 5th-HG signal. As shown in Figure 3c, the $P_{\text{max}}/P_{\text{min}}$ ratio is as high as $\sim$83 with several strong resonance peaks in the range of 2.46–3.30 eV. The highest peak is at $\sim$2.85 eV, close to the C exciton state formed by the band-nesting effect, indicating a strong enhancement of 5th-HG assisted by coupling to the C exciton. The peak at $\sim$3.04 eV is not prominent, so we do not discuss it here. Another strong peak is observed at $\sim$2.6 eV and could arise from other excited exciton states or the electronic band edge, although its origin is still unclear.

Further, we carry out experiments to investigate optically controlled 5th-HG with the control light at 3.1 eV. Figure 3d shows the time-resolved broadband modulation of the 5th-HG signal ($\Delta P_{5\text{th}}$) from 2.80 to 2.92 eV with an average control power $P_c \approx$ 100 nW. The 5th-HG signal can be strongly bleached with the control light, and the highest modulation depth has been found at a generated photon energy of $\sim$2.90 eV, close to the C exciton. To observe the dynamics, we extract the time-resolved modulation of the 5th-HG at a generation photon energy of $\sim$2.85 eV. As shown in Figure 3e, $\Delta P_{5\text{th}}$ behaves similarly to $\Delta P_{4\text{th}}$ first decreasing sharply within 100 fs and then recovering with two different time constants ($t_1 \approx 1.7 \pm 0.04$ ps and $t_2 \approx 143.2 \pm 26$ ps). Figure 3f indicates that the fifth-harmonic modulation $\Delta P_{5\text{th}}$ first decreases linearly when the control power is below $\sim$100 nW, and then tends to be saturated when increasing the control power. The maximum modulation depth can reach up to $\sim$84% with a control light power of $\sim$500 nW ($I_c \sim$ 7.18 GW/cm²).

The underlying mechanism of solid-state HHG is generally described as a combined effect of interband polarization and intraband oscillations. Especially, when the generated harmonic photon energy is larger than the bandgap, the solid-state HHG operating with significant overlap of atomic orbitals gives rise to a probability of interband polarization, which we attribute as the mechanism explaining our experimental results: HHG is highly dependent on the amount of valence carrier states, and thus its strength is modified by interband excitations. This is illustrated by our HHG measurements (Figures 2c and 3c), which support a large enhancement at the C exciton. The observable enhancements of 4th-HG and 5th-HG in amplitude are as high as $\sim$17 and $\sim$83 times compared to the lowest measured nonlinear signal wavelengths. The deviation in the power dependence from the perturbative power-law $P_0^\alpha$ with exponents $\alpha$ of 4 and 5 for 4th-HG- and 5th-HG (see Figures 2b and 3b) further points to a dominant effect played by multiphoton state-related resonances. With increasing seed-light power, the HHG signal tends to be strongly saturated (Figures 2b and 3b), a clear multiphoton saturable absorption effect is observed.

We understand that by illuminating with control light, a redistribution of valence carriers is produced, which affects HHG and, thus, enables its modulation. More precisely, we attribute this effect to the bleaching of valence carriers by the control light (i.e., promotion to the conduction band), which reduces the interband polarization strength and, therefore, decreases the HHG efficiency. This mechanism supported by our time-resolved HHG measurements in Figures 2e and 3e clearly shows that the bleaching process happens very fast within 100 fs. Subsequently, HHG starts recovering with two different time constants along with the relaxation of photocarriers. The fast time constant can be understood as the result of carrier relaxation processes, while the slow time constant should be associated with slower electron–hole recombination processes. Our results for the dependence of the modulation on the power of the control light (Figures 2f and 3f) also confirm that the bleaching of valence carriers with the control light can significantly reduce the HHG intensity. Considering the high modulation depth of 4th-HG (5th-HG) by as much as 95% (84%), we further corroborate that interband polarization contributes dominantly to HHG in the explored energy region.

In our experiment, we also extend our optical control method to HHG in graphene (see Figures S5 and S6 in section 8 of Supporting Information). Compared to the transient responses in MoS2, the decay time of photon carriers in graphene is much faster because of its higher carrier relaxation rate, while the modulation depth in graphene is relatively small due to its low carrier density. This further highlights the importance of the coupling to the exciton states for large modulation in MoS2.

To fully understand these experimental results, we calculate the high-order nonlinear optical response of monolayer MoS2. In the past, the intrinsic complexity of the nonlinear optical response of nanomaterials has stimulated the development of quantitative theoretical prediction capabilities. Indeed, advanced theoretical approaches have been introduced in recent years to simulate nonlinear optical phenomena, such as second-harmonic generation (SHG), and third-harmonic generation (THG), and formulate predictions in good agreement with experimental observations. Here, we use density functional theory in a real-time approach to calculate $ab initio$ the nonlinear optical response of monolayer MoS2. In particular, we use many-body perturbation theory to accurately obtain the electronic structure of the material within the G0W0 approximation and further include excitonic effects by solving the Bethe–Salpeter (BS) equation as implemented in the YAMBO code. Subsequently, we perform real-time simulations to predict the spectral dependence of the nonlinear susceptibilities associated with harmonic generation at second-to-fifth orders.

To understand the nonlinear optical response of the material in detail, we first discuss optical transitions captured by the linear dielectric function. The electronic structure and absorption spectrum of the material (see Supporting Information for parameter sets) are shown in Figure S3. The G0W0 method produces a direct electronic bandgap of $\sim$2.71 eV in monolayer MoS2. In addition, BS calculations for the spectral positions of the optical band-edge excitons $1_s$ and $1_s$ yield 2.04 and 2.18 eV, respectively, while the $1_s$ exciton is also located below the bandgap at $\sim$2.69 eV, and contributions from higher-order excitonic states $\sim$2.41 eV are discernible in the spectrum.

Figure 4a,b shows our simulated spectra for second-to-fifth harmonic generation from monolayer MoS2 over a broad spectral frequency range. It is worth noting that the reduced spatial symmetry of the monolayer enables the generation of even-order harmonics in addition to odd-order ones, in contrast to bulk and even-layered systems of the same crystal. In
particular, the SHG spectrum (blue curve in Figure 4a) is in excellent agreement with previous experimental\textsuperscript{59} and theoretical\textsuperscript{58} results, thus demonstrating the reliability of our calculation methods. Second- and higher-order harmonics in monolayer MoS\textsubscript{2} exhibit multiple resonances associated with excitonic states, whose spectral positions are visible in the absorption spectrum (Figure 4a,b). More precisely, we find that the highest third- and fifth-harmonic signals ($|\chi^{(3)}| \approx 3.8 \times 10^{-18}$ m$^2$/V$^2$ and $|\chi^{(5)}| \approx 6.6 \times 10^{-36}$ m$^4$/V$^4$, respectively) originate in coupling to A and B excitons, while the strongest even-order harmonic responses are stemming from the 1$\text{st}_\text{C}$ exciton ($|\chi^{(2)}| \approx 4.8 \times 10^{-9}$ m/V and $|\chi^{(4)}| \approx 1.1 \times 10^{-27}$ m$^4$/V$^4$). These results agree with previous observations of high even-order harmonic response in the band nesting regions of monolayer TMDs.\textsuperscript{11} To more clearly illustrate this, we calculate the relative output intensities for harmonic generation in resonance with 1$\text{st}_\text{A}$, 1$\text{st}_\text{B}$, and 1$\text{st}_\text{C}$ excitons under the same excitation intensity (see Supporting Information for details), as shown in Figure 4c. While all excitons show comparable output intensities for odd-ordered harmonics, 1$\text{st}_\text{C}$ excitons are found to be more efficient at even orders. In addition, we also study the output intensities observed at resonance with the 1$\text{st}_\text{C}$ exciton (our experimental setup allows us to obtain systematic results only around the 1$\text{st}_\text{C}$-exciton region). While theoretical and experimental output intensities are comparable for SHG, the results differ more as the harmonic order increases, although the qualitative trend is similar in both theory and experiment. Note that different excitation intensities would induce variations in efficiency due to the different order power dependence of the nonlinear optical signals.

To obtain a microscopic understanding of the time-dependent HHG process, we further calculate the electronic-level occupation dynamics in the system based on first-principles parameters and phenomenological decay rates. In particular, we simulate the transition process of charge carriers excited by the control laser field to the ground level by radiative and nonradiative recombination events. The number and distribution of excited electrons in the conduction bands evolve in time according to the rate equation

$$\dot{\Omega}(\varepsilon) = \int_{\varepsilon}^{\infty} \mathrm{d}\varepsilon' \gamma(\varepsilon', \varepsilon) [\rho(\varepsilon') - \Omega(\varepsilon')] \Omega(\varepsilon')$$

$$+ \Gamma_{\text{rad}}(\varepsilon, t) \rho(\varepsilon) - \Omega(\varepsilon) \Omega(\varepsilon - \hbar \omega)$$

$$- \int_{\varepsilon}^{\infty} \mathrm{d}\varepsilon' \gamma(\varepsilon, \varepsilon') [\rho(\varepsilon') - \Omega(\varepsilon')] \Omega(\varepsilon')$$

$$- \Gamma_{\text{nonrad}}(\varepsilon, t) \rho(\varepsilon + \hbar \omega) - \Omega(\varepsilon + \hbar \omega) \Omega(\varepsilon)$$

where $\rho(\varepsilon)$ and $\Omega(\varepsilon)$ denote the density of states (see Figure S4a) and the electron occupation density as a function of electron energy $\varepsilon$, respectively. Also, $\Gamma_{\text{rad}}(\varepsilon, t) = (2\pi \varepsilon E_0^2 / \hbar) e^{-\varepsilon^2 / (\hbar^2 \Delta^2)}$ describes the excitation rate associated with the electric-field amplitude of the external control light $E_0$. We consider excitation by a laser pulse with $\Delta = 150$ fs duration centered around a photon energy $\omega = 3.1$ eV. In $\Gamma_{\text{nonrad}}(\varepsilon, t)$, we introduce $\Delta^2 \epsilon_{\text{excitons}}$ as the Brillouin zone (BZ) average of the squared transition dipole matrix elements between states with energies $\varepsilon$ and $\varepsilon - \hbar \omega$. We consider two types of processes contributing to the recombination rate, radiative and nonradiative, so we have $\gamma = \gamma_{\text{rad}} + \gamma_{\text{nonrad}}$. The nonradiative decay rate $\gamma_{\text{nonrad}}$ is considered to be $1/(100 \text{ fs})$ for transitions with energies less than 0.2 eV as a phenomenological way of incorporating intraband processes assisted by phonon creation. In addition, for the radiative decay rate, we consider $\gamma_{\text{rad}}(\varepsilon, \varepsilon') = \frac{\hbar^2 |\chi(\varepsilon', \varepsilon)|^2}{\hbar^2}$ which describes recombination assisted by the emission of one photon. Unfortunately, this expression for $\gamma_{\text{rad}}(\varepsilon, \varepsilon')$ results in recombination times that are much longer than those observed experimentally ($\sim 30$ ps), thus indicating that additional recombination mechanisms other than photon emission must be dominant and in particular recombination assisted by exciton creation.\textsuperscript{62,63} In this line, Auger electron–hole recombination assisted by bound excitons has been observed in MoS\textsubscript{2},\textsuperscript{64,65} while defect-assisted recombination is known to affect the recombination rates in this material as well.\textsuperscript{66} Assuming a similar role of the transition matrix elements as in radiative recombination, we can phenomenologically account for exciton-mediated recombination by correcting $\gamma_{\text{rad}}(\varepsilon, \varepsilon')$ through a multiplicative factor $3 \times 10^5$, which we determine by comparison with the experimental results. This argument is supported by the high local density of optical states at the layer position, dominated by coupling to excitons, as argued in section 7 of the Supporting Information. The resulting temporal evolution of the charge carrier energy distributions is plotted in Figure 4d, where we observe that short duration after the control laser is applied, the excited holes are grouped within a small energy window at the top of the valence band via fast transitions. In addition, the electron occupation of the conduction bands obtained from this model (dashed-yellow curves in Figures 4e and 4f) are in good correspondence with the experimental observations and the biexponential fit.

In conclusion, we have demonstrated active control of HHG at various orders in materials of atomic thickness. A strong enhancement of HHG (\~83 times) has been determined for monolayer MoS\textsubscript{2} at harmonics resonating with the 1$\text{st}_\text{C}$ excitonic state. By optically modifying the distribution of charge carriers, HHG in monolayer MoS\textsubscript{2} and few-layered graphene has been controlled with high modulation depth (up to 95%), low power consumption (250 pJ), ultrafast speed (several picoseconds), and small footprints (atomic thickness). Further, we have
developed a quantitative theoretical analysis for the nonlinear optical responses associated with second-to-fifth-order processes, as well as the HHG efficiency and its relation to the temporal evolution of the carrier density distribution. Our demonstration of optically controlled HHG paves a new route for engineering extreme nonlinear optics in atomically thin materials with potential application in active optical devices.

■ ASSOCIATED CONTENT

Supporting Information
The Supporting Information is available free of charge at https://pubs.acs.org/10.1021/acs.nanolett.2c02711.

(1) Optical properties of monolayer MoS$_2$; (2) time-resolved THG modulation below the bandgap; (3) computational details; (4) derivation of relative intensities; (5) carrier dynamics; (6) calculation of density of states and dipole matrix elements; (7) enhancement of recombination transition rates by coupling to excitons; (8) optically controlled HHG in few-layer graphene (PDF)

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Author Contributions
Y.W. and F.I. contributed equally to this paper. Y.W. and Z.S. conceived the idea. Y.W. performed the measurements assisted by Y.D., X.H., S.D., Y.Z., and X.B. and S.L. synthesized the MoS$_2$ crystal. Y.W. and X.H. exfoliated and characterized the few-layered graphene. F.J.G.d.A. proposed the theoretical model. F.I. performed the theoretical calculations. All authors discussed the results and commented on the manuscript.

Funding
The authors acknowledge the financial support from Aalto Centre for Quantum Engineering, Academy of Finland (Grants 314810, 333982, 336144, and 336818), the Academy of Finland Flagship Programme (Grant 320167,PREIN), Foundation for Aalto University Science and Technology (Grant 630001), ERC (Advanced Grants 834742-ATOP and 789104-eNANO), Spanish MCINN (Grants PID2020-112625SB-I00 and CEX2019-000910-S), Generalitat de Catalunya (CERCA and GAUR), and Fundació Cellex and Mir-Puig.

Notes
The authors declare no competing financial interest.

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