Excitonic properties of semiconducting monolayer and bilayer MoTe$_2$

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MoTe$_2$ belongs to the semiconducting transition-metal dichalcogenide family with certain properties differing from the other well-studied members (Mo,W)(S,Se)$_2$. The optical band gap is in the near-infrared region, and both monolayers and bilayers may have a direct optical band gap. We first simulate the single-particle band structure of both monolayer and bilayer MoTe$_2$ with density-functional-theory-$GW$ calculations. We find a direct (indirect) electronic band gap for the monolayer (bilayer). By solving in addition the Bethe-Salpeter equation, we find similar energies for the direct exciton transitions in monolayers and bilayers. We then study the optical properties by means of photoluminescence (PL) excitation, reflectivity, time-resolved PL, and power-dependent PL spectroscopy. With differential reflectivity, we find a similar oscillator strength for the optical transition observed in PL in both monolayers and bilayers suggesting a direct transition in both cases. We identify the same energy for the $B$-exciton state in the monolayer and the bilayer. Following circularly polarized excitation, we do not find any exciton polarization for a large range of excitation energies. At low temperatures ($T = 10$ K), we measure similar PL decay times on the order of 4 ps for both monolayer and bilayer excitons with a slightly longer one for the bilayer. Finally, we observe a reduction of the exciton-exciton annihilation contribution to the nonradiative recombination in bilayers.

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I. INTRODUCTION

Monolayers (MLs) of group-VI transition-metal dichalcogenide (TMD) form a new class of semiconducting materials with exciting properties for electronics and optoelectronics applications [1–4]. MoS$_2$ was the first studied material [5,6] quickly followed by extensive studies on binary WS$_2$,MoSe$_2$,WSe$_2$ [7,8], and ternary alloys [9–11]. In the 2H hexagonal structure, these materials all share common properties: (i) an indirect-to-direct band-gap crossover when the material is thinned down to the monolayer limit [5,6] where the direct gap is located at the $K^\pm$ points of the hexagonal Brillouin zone, (ii) strong binding energies for excitons (Coulomb bound electron-hole pairs) of several hundreds of meV [12–14], and (iii) spin- and valley-dependent optical selection rules due to the lack of crystal inversion symmetry and large spin-orbit coupling (SOC) [15–19]. Nevertheless, several differences in the physical properties between semiconducting TMD MLs materials can be pointed out. The amplitude of the valence spin-orbit splitting varies from less than 200 meV for Mo(S,Se)$_2$ MLs to more than 400 meV for W(S,Se)$_2$ MLs. The sign of the splitting between intravalley bright and dark excitons is the opposite in MoSe$_2$ and WSe$_2$ MLs [20], dramatically affecting their optical properties [11,21,22]. Finally, despite comparable structural and optical qualities, the valley-/spin-polarization properties probed in optical spectroscopy are very different between materials. For MoS$_2$ MLs, the degree of circular polarization of the photoluminescence (PL) decreases monotonously when the detuning of the excitation laser energy relative to the excitonic ground state increases [23,24]. For WSe$_2$ and WS$_2$ MLs, circularly polarized PL has been reported even for excitation energies far from the resonance, but the polarization degree is exalted for excitation energies in resonance with exciton excited states (2$s$ or 2$p$) [14,25]. In contrast, no PL polarization has been observed in MoSe$_2$ MLs except for quasiresonant excitation [24].

Recently, the family of semiconducting TMD monolayers expanded with a fifth binary material: MoTe$_2$ [26,27]. Two striking properties have been highlighted by the first studies on this material: The optical band gap of MoTe$_2$ MLs is in the near infrared (IR) (1.1 eV at room temperature [26]) whereas it lies in the red part of the visible spectrum for S- and Se-based compounds. Second, the luminescence yield of MoTe$_2$ bilayers (BLs) is on the same order of magnitude as the luminescence yield of the ML opening a debate on a possible direct band gap for BLs [27,28]. Nevertheless, the high luminescence yield of BLs is up to now the only argument pointing towards a direct band gap. In addition, many properties, such as radiative lifetime, spin/valley polarization, and the energy of exciton excited states, have not been measured in this material yet. In this paper we partially fill this gap by studying theoretically and experimentally the optical properties of both MoTe$_2$ MLs and BLs. The paper is organized as follows. In the next section, we use DFT + $GW$ calculations to simulate the band structures of both the ML and the BL in a single-particle picture (no excitonic effects). Remarkably we find an indirect electronic band gap for the BL. We then solve the Bethe-Salpeter equation (BSE) to include the strong electron-hole Coulomb attraction when calculating the energy and the oscillator strength of direct excitonic transitions. Section III is dedicated to stationary PL measurements. Experimentally, we measure the splitting between $A$ and $B$ excitons by PL excitation spectroscopy.
and reflectivity and compare the results with $G_0W_0 + $ BSE calculations. We do not find any Stokes shift between the PL of the A exciton and the strong signature observed in reflectivity in both the ML and the BL indicates a direct optical band gap in both systems. We also do not detect any measurable PL polarization for excitation energy as close as 60 meV above the energy of the A exciton suggesting a behavior similar to MoSe$_2$ MLs. In Sec. IV, we use time-resolved PL to measure the exciton lifetime in the range of a few picoseconds in both MLs and BLs with a slightly longer one for BLs. Finally, in Sec. V, we show that exciton-exciton annihilation (EEA) is larger in MLs than in BLs.

II. ELECTRONIC BAND-STRUCTURE CALCULATION

In recent years, DFT + $GW$ methods have been successfully applied to calculate the electronic band gap of two-dimensional (2D) semiconductor materials ($E_G$), see, for instance, Refs. [14,29–32]. In the present paper, the exploration of the electronic structure and optical properties of ML and BL MoTe$_2$ has been performed using the VASP code [33,34]. It uses the plane-augmented-wave scheme [35,36] to treat core electrons when 14 electrons for Mo and 6 for Te ones are explicitly included in the valence states with a plane-wave energy cutoff of 400 eV. The Heyd-Scuseria-Ernzerhof hybrid functional [37–39] is used to build the needed wave functions, based on 600 electronic states, to calculate the full-frequency-dependent quasiparticle band structure at the $G_0W_0$ level of theory [40], including SOC but not the excitonic effects at this stage. This particular choice of the computational settings has been discussed in detail in Ref. [20]. A grid of $(12 \times 12 \times 1)$ $k$ points has been used in conjunction with a vacuum height of at least 17 Å for both the ML and the BL systems. For the latter, we have selected the stacking geometry of $AA'$ type with the optimized interlayer distance is 7.00 Å when van der Waals (VdW) forces are taken into account via the optB86b-VdW scheme [42].

In Fig. 1(a), the resulting DFT + $G_0W_0$ band structure of the MoTe$_2$ ML is shown after a Wannier interpolation procedure performed by the WANNIER90 program [43]. Its main features agree well with previous theoretical studies [29,44]. It has a direct electronic band gap in the $K$ valley with a value of 1.72 eV, a SOC splitting in the valence band of 275 meV whereas it is −58 meV in the conduction band. The negative sign for the conduction band SOC means that both conduction-band minimum and valence band maximum have the same spin, i.e., the lowest-lying interband transition is bright. In comparison the BL band structure at the same level of theory is given in Fig. 1(b). If the direct $K_v$-$K_c$ gap remains almost the same with a value of 1.66 eV for the BL, now it appears that the indirect $K_v$-$\Lambda_c$ quasiparticle band gap is the lowest one with a value of 1.60 eV. Thus the interlayer interaction leads to a transition from a direct to an indirect band gap similar to other group-VI TMDs [41] with an energy separation between $K_c$ and $\Lambda_c$ remaining small compared to other TMD systems. Nevertheless, contrary to MoS$_2$ BLs, the valence-band minimum remains in the $K$ valley [41]. The interlayer interaction also tends to enhance the energy separation between the two highest valence states of different spins for each layer, whereas the opposite is true for the lowest unoccupied ones. As a result, the SOC splittings become 304 and −46 meV for the valence and conduction bands, respectively.

With DFT + $GW$ calculations we find that the electronic band gap is direct for the ML and indirect for the BL. But in TMD materials, the optical properties are not governed by the band-to-band recombination. In reality, due to the very large exciton binding energy ($E_B \sim 500$ meV), the PL spectrum is dominated by the ground exciton transition (also called optical band-gap $E_{PL} = E_G - E_B$). To calculate the optical band gap,
we need to include the electron-hole Coulomb interaction into the model by solving the BSE. Practically BSE spectra are obtained in the Tamm-Dancoff approximation by using the six highest valence bands and the eight lowest conduction bands for the MoTe

2, which reduces significantly the binding energy at (i) of a weaker intralayer screening and (ii) of a larger interlayer distance when passing from S composition. We also observe this trend for exciton binding energy values. We have obtained a binding energy of around 0.16 eV, in good agreement with an experimental determination reported in Ref. [47]. So the ratio between bulk and ML exciton binding energies is around 33%, whereas for MoS

2 it is on the order of 20% if one considers roughly the exciton ML binding energy around 0.5 and 0.1 eV for the bulk.

From the ML spectrum we can also identify a small shoulder at 0.04 eV below the $X_A^I$ peak, which can be safely assigned to a transition associated with the 2$s$ state of the A exciton. Indeed this transition appears as soon as the $X_A^I$ peak has non-negligible oscillator strength and is of interlayer character. Indeed it involves the fifth and sixth conduction bands which clearly possess a delocalized character over the two layers and is composed mainly of $d_{xz}$ and $d_{yz}$ orbitals of the two Mo atoms. Note that this kind of weak interlayer transition in the vicinity of the A peak has been previously reported not only in the case of the MoS

2/WS

2 heterostructure [30], but also for bilayers [41].

III. CONTINUOUS-WAVE SPECTROSCOPY

Our DFT calculations predict at the GW level a direct gap for the ML and an indirect gap for the BL with a direct gap only 60 meV above in energy. This energy difference is very small once excitonic effects ($\sim$500 meV) are taken into account. This motivates the optical spectroscopy studies described in this section where ML and BL emission energies and intensities can be determined.

For experimental studies of optical properties, we use MoTe

2 flakes obtained by micromechanical cleavage of a bulk MoTe

2 crystal (supplied by the company 2D semiconductors) on a 90-nm SiO

2/Si substrate using viscoelastic stamping [48]. For reflectivity measurements, we use transparent quartz substrates. The ML and BL regions are identified by optical contrast [see the inset of Fig. 2(a)] and very clearly in PL

FIG. 2. (a) PL spectra of the MoTe

2 ML (black) and BL (red) at 10 K. The inset: optical microscope image used for the identification of ML and BL regions. (b) PL spectra of the MoTe

2 ML from 10 to 100 K. The sample is excited with a cw HeNe laser (633 nm) at a power of 50 $\mu$W.
the variation of the presents the PL excitation (PLE) spectra corresponding to excited excitonic states is of particular importance. Figure 3(a) on the low-energy part of the spectrum and between temperatures we obtained on chemically treated MoS$_2$ MLs is [51,52]. In addition to the two main peaks we observe features the Coulomb interaction, and the spin-orbit coupling properties indicates that disorder may impact the radiative recombination, optical quality as high as for the best MSe$_2$-(M$^+$-$\Delta$M$^{-}$) peaks of $\Delta 1$ and trion $T$ (at 1.159 eV for a ML) for the A transition. The energies of the peaks are in good agreement with the previously reported PL spectra of MoTe$_2$ MLs and BLs [26–28]. For the MoTe$_2$ ML, we measure a separation between the $X_{A}^{1s}$ and the $T$ peaks of 25 meV in agreement with the measurement of the binding energies of positively and negatively charged excitons in a field effect structure [44] (24 and 27 meV, respectively). For the BL, the separation between peaks is smaller (18 meV), but we cannot unambiguously attribute the low-energy peak to the trion signature as no charge tunable device based on BL has been reported yet. The full width at half maximum is 7 meV for both $X_{A}^{1s}$ and $T$ in the ML indicating an optical quality as high as for the best MSe$_2$-(M$^+$-$\Delta$M$^{-}$) ML [49] samples. For comparison, the smallest linewidth at low temperatures we obtained on chemically treated MoS$_2$ MLs is typically 15 meV [50]. Thus, MoTe$_2$ is particularly suitable to study the complex exciton/trion fine structure of TMD MLs. Nevertheless, the measured linewidth is still broader than the expected homogeneous linewidth (see also Sec. IV). This indicates that disorder may impact the radiative recombination, the Coulomb interaction, and the spin-orbit coupling properties [51,52]. In addition to the two main peaks we observe features on the low-energy part of the spectrum and between $X_{A}^{1s}$ and $T$ peaks. We attribute them to complex localized states [marked as Loc in Fig. 2(a)] as their contribution to the PL spectrum vanishes when the temperature increases [see Fig. 2(b)]. We also want to point out that the amplitude of these localized states varies from a ML-to-ML sample making the origin of these features difficult to attribute at this stage.

Several important characteristics of TMD MLs (optical generation of valley polarization, second-harmonic generation efficiency) are known to strongly depend on the laser excitation energy as the light-matter interaction is strongly enhanced at the excitonic resonances [14,24,53,54]. Therefore, probing the excited excitonic states is of particular importance. Figure 3(a) presents the PL excitation (PLE) spectra corresponding to the variation of the $X_{A}^{1s}$ PL intensity as a function of the laser energy. The PLE spectra depend on the absorption at the excitation energy and how the carriers relax down to the $X_{A}^{1s}$ state. A clear resonance of the $X_{A}^{1s}$ PL peak is observed at 1.459 eV for both the ML and the BL which corresponds to the signature of the ground B exciton state $X_{B}^{1s}$. The energy of the $X_{B}^{1s}$ state is also confirmed in the reflectivity spectrum of Fig. 3(b), which scales like the absorption spectrum without being sensitive to carrier relaxation as the...
PLE measurements. We thus find a splitting \( B-A \) of 270 meV for the ML, in agreement with the room-temperature reflectivity measurement of Ruppert et al. (260 meV) [26]. This is also in good agreement with the BSE spectrum of Fig. 1(c) where we calculate a splitting of 330 meV. We notice that this splitting is larger than for the other Mo-based TMDs. Finding exactly the same \( X^{1s}_B \) resonance energy for both a ML and a BL (while there is a 35-meV shift for the \( X^{1s}_A \) resonance energy) is also interesting and in good agreement with the calculations of Figs. 1(c) and 1(d).

Contrary to measurements on WSe\(_2\) MLs [55], we do not find a clear signature of the \( X^{2a}_A \) excited state in the PLE spectrum of the MoTe\(_2\) ML. A possible explanation is that this state is very close to the \( X^{1s}_B \) state, such as in MoSe\(_2\) MLs [56] and cannot be distinguished in the PLE spectrum. This is supported by the \( GW\)-BSE calculations of Fig. 1(c) where we find a separation \( X^{1s}_B-X^{2a}_A \) of 40 meV thus smaller than the width of the PLE peak in Fig. 3(a). Two-photon experiments would be an efficient way to probe the \( X^{2a}_A \) state without being sensitive to the \( X^{1s}_B \) state [56]. However this requires exciting the ML with energies lower than 0.75 eV, which is beyond the tuning range of conventional laser systems. For the BL we can point out that the resonance peak at 1.45 eV observed in both PLE and reflectivity spectra is clearly broader than in the ML case and could be a sign of the \( X^{2a}_A \) contribution.

Using the ratio of integrated BL PL versus ML PL intensity is not sufficient to distinguish between direct and indirect optical transitions. Indeed, the laser energy \( E_{\text{laser}} \) used for excitation plays an important role as can be seen in Fig. 3(c). For example, the BL PL emission for \( E_{\text{laser}}=1.319 \text{ eV} \) is more intense than the ML PL. A stronger argument is given by the reflectivity spectrum of Fig. 3(b). Because the flakes are transferred onto a thick transparent quartz substrate in this experiment, the differential reflectivity given by \( \Delta R/R = \frac{(R_{\text{flakes}} - R_{\text{quartz}})}{R_{\text{quartz}}} \) qualitatively resembles the absorption spectrum [57–59]. For both ML and BL samples, a strong resonance in \( \Delta R/R \) is observed at the exact energy of the \( X^{1s}_B \) PL peak. Remarkably, the strength of this resonance is similar in MLs and BLs. These results are a strong indication that the BL PL transition corresponds to a direct optical gap at this energy.

TMD MLs obey chiral optical interband selection rules [15] that allow for optical excitation in either the \( K^+ \) or the \( K^- \) valley depending on excitation laser helicity. For the range of excitation energies shown in Fig. 3(a), we do not measure any significant circular (linear) polarization following circularly (linearly) polarized excitation. This behavior is in agreement with the very recent zero polarization reflectivity measured at a zero magnetic field [60]. This is also similar to what was observed for the MoSe\(_2\) ML [23]. Only one recent study reports a circular PL polarization degree of 20\% in the MoSe\(_2\) ML for a difference between the laser energy and the energy of the \( X^{1s}_A \) emission of \( \Delta E = E_{\text{laser}} - E[X^{1s}_A] = 60 \text{ meV} \) [24]. Longitudinal acoustic (LA) phonon-assisted intervalley scattering was proposed as the cause of depolarization in MoSe\(_2\) and MoS\(_2\) MLs [24]. The larger value of the LA phonon for MoSe\(_2\) (30 meV) [61] as compared to MoSe\(_2\) (19 meV) [62] would explain why polarization can be observed in MoS\(_2\) for larger \( \Delta E \)'s. Interestingly, the LA phonon energy of MoTe\(_2\) is even smaller (12 meV) [63] than that of MoSe\(_2\). Thus, only excitation very close to the resonance may initiate valley polarization in this material. Unfortunately, for \( \Delta E < 60 \text{ meV} \), a strong Raman-scattering signal is superimposed on the PL signal resulting in unreliable polarization measurements.

**IV. TIME-RESOLVED MEASUREMENTS**

The exciton binding energy in a MoTe\(_2\) ML can be roughly estimated by taking the difference between the electronic band gap provided by the DFT-\( GW\) calculation [1.72 eV, see Fig. 1(a)] and the optical band gap measured in PL [1.18 eV, see Fig. 2(a)]. We find a binding energy of \( \sim 540 \text{ meV} \) in agreement with previous estimation of Yang et al. [44]. Such a strong exciton binding energy and the associated strong oscillator strength are very interesting properties for strong light-matter coupling studies in ML TMDs [64,65]. Consequently, a short PL emission time can be expected, and measuring the exciton radiative lifetime is thus crucial. TRPL is the ideal spectroscopy tool for such a measurement. In our previous works, we measured low-temperature PL decay times of 4 ps for Mo\(_2\) MLs [54] and 2 ps for Mo\(_2\) and WSe\(_2\) MLs [66]. Our temperature-dependent study on a MoSe\(_2\) ML suggested that this decay time corresponds to the radiative lifetime of excitons before thermalization occurs. In Fig. 4, we present the low-temperature PL dynamics of \( X^{1s}_B \) for both Mo\(_2\) MLs and BLs. The excitation wavelength of 850 nm (1.458 eV) is chosen to match with the enhanced absorption at the \( X^{1s}_B \) exciton transition [see Fig. 3(a)]. For the ML, we observe a monoexponential decay that can be fitted with a characteristic time of \( \tau = 3.4 \pm 0.5 \text{ ps} \). This dynamic is clearly longer than the time resolution of our setup measured by detecting the laser pulse backscattered from the sample on the streak camera (shaded area). Following Ref. [66], we interpret this fast decay time as the radiative lifetime of the excitons. A key argument is that the \( X^{1s}_A \) PL intensity is constant in the range of 10–40 K [see Fig. 2(b)] excluding any role of nonradiative channels at low temperatures. In the context of the optical generation of valley polarization, the short PL emission time \( \tau_{ML} \) of 3.4 ps implies that either the valley depolarization...
time is considerably shorter than 3.4 ps or the initially created valley polarization is negligible. For the BL, the decay is biexponential. We attribute the longest decay time (23 ps) to the presence of localized states at the same energy as the free exciton peak. The shortest decay time $\tau_{\text{BL}} \sim 4.3 \pm 0.5$ ps can be attributed to the lifetime of free excitons in the BL. It is thus very similar to $\tau_{\text{ML}}$ and six times faster than for the indirect transition in a WSe$_2$ BL [67]. This suggests a direct transition in a MoTe$_2$ BL in agreement with the reflectivity spectra of Fig. 3(b). Nevertheless, we want to point out that contrary to the ML, the PL intensity of the BL decreases as soon as the temperature is raised. We thus cannot exclude that this decay time is also governed by nonradiative recombination. Interestingly, this time is slightly longer than $\tau_{\text{ML}}$, which might hint at the fact that the radiative lifetime of excitons in the BL is longer than the radiative lifetime of excitons in the ML. Several hypotheses can be proposed at this stage. Even if the PL of a MoTe$_2$ BL originates from a direct transition, it was reported that the optical band gap of a MoTe$_2$ BL may be borderline direct/indirect [27,28]. It is thus not surprising to observe a smaller rate of radiative recombination in the BL. Second, potential fluctuations with a high spatial frequency have been proposed to explain the significant discrepancy between the measured exciton lifetime (a few picoseconds) and the theoretical radiative lifetime of free excitons (a few hundreds of femtoseconds) [66]. We expect larger Bohr radius excitons in the BLs to be more sensitive to these fluctuations than smaller Bohr radius excitons in the ML and consequently to yield longer radiative lifetimes.

V. EXCITON-EXCITON ANNIHILATION

Nonradiative recombination channels are known to play a major role in the poor luminescence yield measured at room temperature in TMD MLs [68]. In addition to defect-related recombination, EEA is known to be very efficient in MoS$_2$, MoSe$_2$, WSe$_2$, and WS$_2$ at room temperature even for moderate excitation power density [68,69]. In Fig. 5, we plot the variations of the $X_{A,1}^{\text{PL}}$ intensity as a function of excitation power for both the ML and the BL at 10 and 200 K. For the ML, the $X_{A,1}^{\text{PL}}$ intensity scales linearly with $10^5$ K whereas it scales sublinearly ($\sim P^{0.8}$) at 200 K for excitation powers larger than 10 $\mu$W. This might be a consequence of the thermal activation of EEA processes. At 10 K, the exciton diffusion is too small, and the radiative lifetime is so short that EEA does not compete with radiative recombination. In contrast, when the temperature increases, the mobility of the excitons increases [69,70] and so does the effective radiative lifetime of the excitons due to thermalization [22,71,72]. This enhances the sensitivity to many-body interactions. Remarkably, for the BL, the situation is different. As shown in Fig. 5, the $X_{A,1}^{\text{PL}}$ intensity increases linearly with the excitation power for both temperatures suggesting a reduced EEA rate as compared to the ML. We also observe the linearity at room temperature (not shown here). Such a property combined with an efficient optical transition would make the MoTe$_2$ BL a very promising candidate for optoelectronics applications requiring high carrier densities including laser or concentrating solar cells [73]. Unfortunately, a clear conclusion cannot be drawn at this stage. First, we notice that, at 10 K, the PL intensities of the MLs and BLs are on the same order of magnitude whereas at 200 K, the BL intensity is one order of magnitude lower than the ML intensity (for the same excitation power). This could be due to a higher defect density in the BL. Passivation capping or chemical treatments [68,74,75] would help to study the real influence of defects on the optical properties of the MoTe$_2$ ML and BL. A second explanation would be that the indirect transition plays a more important role at elevated temperatures (due to thermal activation or direct-to-indirect crossover). Actually, Yuan and Huang already measured the EEA rates in the WS$_2$ ML and BL [76] and found a rate for the BL two orders of magnitude lower than in the ML due to the reduced phonon-assisted EEA of indirect excitons in the indirect band gap WS$_2$ BL. Thus, we can expect a similar behavior if the MoTe$_2$ BL is indirect at elevated temperatures. Finally, we cannot exclude that the mobility of excitons is reduced in the MoTe$_2$ BL, which could explain why the EEA processes are not visible even at high temperatures.

In conclusion, we studied the optical properties of the MoTe$_2$ ML and BL. We performed DFT-GW calculations and found a direct electronic band gap for the ML and an indirect one for the BL. With one-photon PLE and reflectivity, we found that the energy of the $B$-exciton state is the same in the ML and the BL in agreement with the BSE calculations. We did not find a clear signature of the $X_{A}^{\text{PL}}$ exciton excited state which may lie close to the $X_{B}^{\text{PL}}$ resonance. We did not detect any circular or linear PL polarization for laser energy as close as 60 meV above the energy of the $X_{A}^{\text{PL}}$ exciton. We measured the exciton lifetimes of the ML and the BL at low temperatures. The lifetime in the BL is slightly longer than the radiative lifetime in the ML but remains significantly faster than for the BLs of other TMD materials. Combined with the observation of similar oscillator strength for $X_{A}^{\text{PL}}$ in the ML and the BL in reflectivity measurements, we can reasonably argue that the low-temperature PL originates from direct transitions in both...
MLs and BLs. Finally, we discussed the observed reduction of the EEA contribution to the nonradiative recombination in the MoTe₂ BL.

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