Enhanced valley-resolved thermoelectric transport in a magnetic silicene superlattice

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Abstract
Electrons in two-dimensional crystals with a honeycomb lattice structure possess a valley degree of freedom in addition to charge and spin, which has revived the field of valleytronics. In this work we investigate the valley-resolved thermoelectric transport through a magnetic silicene superlattice. Since spin is coupled to the valley, this device allows a coexistence of the insulating transmission gap of one valley and the metallic resonant band of the other, resulting in a strong valley polarization $P_v$. $P_v$ oscillates with the barrier strength $V$ with its magnitude greatly enhanced by the superlattice structure. In addition, a controllable fully valley polarized transport and an on/off switching effect in the conductance spectra are obtained. Furthermore, the spin- and valley-dependent thermopowers can be controlled by $V$, the on-site potential difference between A and B sublattices and Fermi energy, and enhanced by the superlattice structure. Enhanced valley-resolved thermoelectric transport and its control by means of gate voltages make the magnetic silicene superlattice attractive in valleytronics applications.

1. Introduction
In addition to manipulating the charge or spin of electrons, another way to control electric current in two-dimensional (2D) materials with a honeycomb lattice structure is by using the valley degree of freedom (DOF) [1–3]. This leads to valleytronics, in which information is encoded by the valley quantum number of the electron [1–12]. A fundamental goal in valleytronics is to generate and detect a controllable valley-polarized current. The seminal proposal of valley filters [4] relies on perfect graphene zigzag nanoribbons, which would harden its experimental realization. Several schemes of valley filters have been proposed based on bulk graphene, which utilize the valley dependence of trigonal band warping [13, 14], strain-induced pseudomagnetic fields [15, 16], and line defects [17]. However, the presence of inversion symmetry in the crystal structure of pristine graphene makes control of the valley DOF difficult. In contrast, silicene that possesses a staggered honeycomb lattice structure is inversion-asymmetric. Although the valley- and spin-polarized currents have been investigated in several works in simple normal/ferromagnetic junctions with one or two barriers [18–20], none of the previous works focuses on the valley-polarized transport of 2D magnetic silicene superlattices. It is well known that superlattices possess many interesting electronic transport properties and band structures, which are mainly determined by the periodicity of the barrier potential rather than by the properties of the individual potential barrier. Thus we expect that a 2D silicene monolayer under a periodic magnetic modulation may be a promising candidate to invent valleytronic devices.

Recently, spin caloritronics [21], the combination of thermoelectrics and spintronics, has received much attention and caused the renaissance of thermoelectricity in spintronic devices. Spin caloritronics covers physical phenomena [22] that are classified as independent electron (such as spin-dependent Seebeck), collective (such as spin Seebeck), and relativistic (such as spin–Hall) effects. Spin caloritronics has been studied extensively in graphene-based devices [23], and a giant charge thermoelectric coefficient was reported in a graphene superlattice [24]. However, the effect of the valley DOF on spin caloritronics in a 2D silicene superlattice has not yet been considered.
In this work we study the valley-resolved thermoelectric transport through a 2D magnetic silicene superlattice by using the transfer matrix method. Since spin is coupled to the valley, we observe the valley and spin-resolved transmission, leading to a strong valley polarization $P_{v}$. $P_{v}$ oscillates with the barrier strength, with its magnitude greatly enhanced by the superlattice structure. In addition, a controllable fully valley-polarized transport and an on/off switching effect in the conductance spectra are obtained. Furthermore, the spin- and valley-dependent thermopowers can be controlled by the barrier strength, the on-site potential difference between $A$ and $B$ sublattices and Fermi energy, and enhanced by the superlattice structure.

2. Model and Formulation

The geometry of a magnetic superlattice is a 2D silicene under a periodic magnetic modulation (figure 1(a)). Effectively, the system can be regarded as consisting of normal silicene (NS) with thickness $L$ and ferromagnetic silicene (FS) with thickness $D$ arranged alternately, finally connected to two semi-infinite NS electrodes. The coordinate of the $j$th interface between NS and FS is marked by $l_j$. Only a parallel magnetization configuration will be considered. The potential profile of the system is a multiple step-like quantum well structure with $V=\sigma h$ for $l_{j-1} < x < l_j$ ($j = 1, 2, \ldots, N$), with the barrier strength $V$ controlled by the gate voltage; $V=0$ otherwise. $h$ stands for the ferromagnetic exchange field, which can be induced by a proximity effect between a ferromagnetic insulator and silicene, as proposed and realized for graphene [25, 26]. If the superlattice has $N$ periods, it contains $N-1$ NS and $N$ FS, finally coupled to two NS electrodes. The low-energy effective Hamiltonian simply reads [18, 20, 27, 28]

$$H_{\eta,\sigma} = v_F \left( \hat{p}_x \tau_x - \eta \hat{p}_y \tau_y \right) - \Delta_{\eta,\sigma} \tau_z + V$$

(1)

with $\hat{p}_i = -i\hbar \delta_{i,j}$ and $\Delta_{\eta,\sigma} = \eta \sigma \delta_{\omega} - \Delta_\zeta$, $v_F$ is the Fermi velocity, and $\tau = (\tau_x, \tau_y, \tau_z)$ are Pauli matrices in pseudospin space. $\eta = \pm 1$ corresponds to the $K$ and $K'$ valleys, and $\sigma = \pm 1$ denotes the spin indices. $\delta_{\omega}$ stands for the spin–orbit coupling. $\Delta_\zeta$ is the on-site potential difference between $A$ and $B$ sublattices and can be efficiently tuned by a perpendicular electric field applied perpendicular to the plane. The eigenvalues of the Hamiltonian are given by

$$E_{\eta,\sigma} = \pm \sqrt{\left(\hbar v_F k_\perp^2 \right)^2 + \left(\Delta_{\eta,\sigma} \right)^2} + V,$$

(2)

where $k = \sqrt{k_x^2 + k_y^2}$. The bandgap is located at the $K$ and $K'$ points and given by $2 |\Delta_{\eta,\sigma}|$. Thus, different gaps can arise in different spins and valleys. The energy profile of electrons in the superlattice is illustrated in figure 1(b). In the NS we set $V=0$, so from equation (2) one finds that $E_{\eta,\sigma}$ and $E_{\eta,-\sigma}$ (or $E_{\eta,\sigma}$ and $E_{\eta,-\sigma}$) are degenerate, while in the FS the degeneracy is broken. Due to the coupling between the valley and spin DOF, it is possible to simultaneously manipulate the valley and spin DOF in silicene.

With the general solutions of equation (1), the wave functions in the NS and FS are respectively given by

$$\psi_{\eta,\sigma} = \begin{pmatrix} a_{\eta,\sigma} \left( \frac{\hbar v_F k_\perp}{E_{\eta}} \right) e^{ik_x x} + b_{\eta,\sigma} \left( -\frac{\hbar v_F k_\perp}{E_{\eta}} \right) e^{-ik_x x} \\ \varphi_{\eta,\sigma} e^{-ik_y y} \end{pmatrix},$$

(3)
and

$$\Psi_{FS} = \left[ e_{\eta,\sigma} \left( \frac{\hbar v_F k^+}{E_p} \right) e^{i k^+_x x} + d_{\eta,\sigma} \left( -\frac{\hbar v_F k^-}{E_p} \right) e^{-i k^-_x x} \right] e^{i k_y y}$$

(4)

where $k^+_x = k^+_y \pm i \eta k_y$, $E_N = E + \Delta_{\eta,\sigma}$, and $E_p = E + \Delta_{\eta,\sigma} - V_\alpha$. The wave vectors are

$$k_x = \sqrt{E^2 - \Delta_{\eta,\sigma}^2 \cos \alpha / \hbar v_F}, k_y = \sqrt{(E - v_F^2) - \Delta_{\eta,\sigma}^2} + \sqrt{\Delta_{\eta,\sigma}^2 / \hbar v_F},$$

and $k_y = \sqrt{E^2 - \Delta_{\eta,\sigma}^2 \sin \alpha / \hbar v_F}$ with $\alpha$ the incident angle. For the electron’s transmission through the superlattice, consisting of $N$ potential barriers via the wave function continuity at the NS/FS interfaces, we get the relationship of

$$\left( \frac{t_{\eta,\alpha}}{0} \right) = \left[ R_{\eta,0} (d) \right] \left[ R_{\eta,0} (d) T (0) \right] N \left( \frac{1}{\eta,\sigma} \right)$$

(5)

with $T (0) = R_{\eta,0} (d) M_{\eta,0} M_{\eta,0} R_{\eta,0} (d) M_{\eta,0} M_{\eta,0}$ and $d = D + L$, where

$$M_{\eta,0} = \left( \frac{\hbar v_F k^+_+ - \hbar v_F k^-_-}{E_N} \right), M_{\eta,0} = \left( \frac{\hbar v_F k^+_+ - \hbar v_F k^-_-}{E_N} \right), R_{\eta,0} (x) = \left( \begin{array}{cc} e^{ik^+_x x} & 0 \\ 0 & e^{-ik^-_x x} \end{array} \right)$$

After the transmission probability for spin $\sigma$ and valley $\eta$ electrons $T_{\eta,\sigma} = \left| t_{\eta,\sigma} \right|^2$ is derived from the above equation, the current can be written as

$$I_{\eta,\sigma} = \frac{e}{h} \int_{-\infty}^{+\infty} dE N (E) \int_{-\pi/2}^{\pi/2} da \cos \alpha T_{\eta,\sigma} \left[ f^f_L (E) - f^f_R (E) \right]$$

(6)

where $N (E) = \sqrt{E^2 - \Delta_{\eta,\sigma}^2 - W / \hbar v_F}$ with $W$ the width of the silicene sheet is the carrier density of states (DOS) of the left NS electrode. $f^f_\beta (E) = \left[ 1 + \exp \left( \frac{(E - E_F)/k_B T_\beta}{1} \right) \right]^{-1}$ ($\beta = L, R$) stands for the Fermi distribution function in the $\beta$ electrode. The conductance at zero temperature can be obtained as

$$G_{\eta,\sigma} (E) = \frac{e^2}{h} \int_{-\pi/2}^{\pi/2} d\alpha T_{\eta,\sigma} \cos \alpha d\alpha.$$

(7)

By using equation (7), we can define the valley and spin polarization $P_\eta$ and $P_\sigma$: $P_\eta = (G_K - G_{K'}) / (G_K + G_{K'})$ and $P_\sigma = (G_{\eta} - G_{\eta'}) / (G_{\eta} + G_{\eta'})$, where $G_{\eta} = G_{\eta,\sigma} + G_{K,\sigma}$ is the conductance in the spin-$\sigma$ channel, while $G_{\eta} = G_{\eta,\sigma} + G_{K,\sigma}$ represents the conductance in the $\eta$ valley.

In the linear response regime, i.e., $T_L \approx T_R = T$, we calculate the valley- and spin-resolved thermopower $S_{\eta,\sigma}$ [27, 30]:

$$S_{\eta,\sigma} = -\frac{1}{eT} \frac{L_{\eta,\sigma}}{L_{\eta,0,\sigma}}$$

(8)

with $L_{\eta,0,\sigma} = \frac{1}{T} \int dE \left( E - E_p \right)^2 N (E) \int da \cos \alpha T_{\eta,\sigma} \left[ -\theta f (E) \right]$. Thus, one can introduce the spin-resolved thermopower $S_{\sigma}$ with $S_\sigma = S_{K,\sigma} + S_{K',\sigma}$. The charge- and spin-dependent thermopowers $S_c$ and $S_s$ are calculated as

$$S_c = \frac{1}{2} \left( S_T + S_L \right)$$

(9)

and

$$S_s = S_T - S_L$$

(10)

In analogy with $S_\eta$, we also define the valley-dependent thermopower $S_\eta$

$$S_\eta = S_K - S_{K'}$$

(11)

with $S_\eta = S_{\eta,1} + S_{\eta,1}$. This means that like the spin-dependent Seebeck effect, we can use this valley-dependent Seebeck effect to generate a valley voltage bias [27, 30].

**3. Results and discussions**

In figure 2, the dependence of the valley- and spin-resolved transmission probability $T_{\eta,\sigma}$ on $V$ is presented. Here the parameters are $N = 5$, $D = L = 50$ nm, $\lambda_{\eta,\sigma} = 3.9$ meV, $E = h = 5$ meV, $T = 0$ K, $\Delta_{K} = 3.9$ meV in the NS, and $\Delta_{K'} = 11.7$ meV in the FS. For the incident energy $E$ taken here, the transmission of spin-up (spin-down) electrons in the $K'$ ($K$) valley disappears because of no propagating modes at the left electrode, and only spin-up (spin-down) electrons in the $K$ ($K'$) valley are allowed to propagate. $T_{\eta,\sigma}$ and $T_{\eta,\sigma}$ oscillate with $V$, and for a large incident angle $\alpha$ ($\alpha = \pi/6$ and $\alpha = \pi/3$) the transmission magnitudes never decay with increasing $V$. These results are similar to those observed for topological insulator ultrathin films [31] but quite different from those for conventional 2D electron gas. It should be pointed out that the resonant bands are formed in the transmission spectra, which are separated by nonresonant gaps. For a finite superlattice with five potential barriers the transmission probability shows the four-fold resonance splitting for each resonant band [32]. For a
large $\alpha$, the transmission tends to be zero in the gaps. Furthermore, as $\alpha$ increases, each resonant band of both $T_{+}$ and $T_{-}$ narrows, while nonresonant gap is broadened. It is very interesting to note that tunneling through the superlattice shows strong valley-dependent features. We observe a coexistence of the insulating transmission gap of one valley and the metallic resonant band of the other. In this case, a strong valley polarization can be generated, which will be shown in figure 3.

The valley polarization $P_{v}$ as a function of $V$ with different periods $N$ is studied in figure 3(a). As indicated in figure 3(a) $N$ changes from 2 to 14 in steps of two. (b) Valley-resolved conductances $G_{++,}$ and $G_{--,}$ (in units of $G_{0} = e^{2}N (E)/h$) with $N = 10$ versus $V$. Other parameters are the same as in figure 2.

Figure 2. Valley-resolved transmission coefficients as a function of the barrier strength $V$ modulated by the gate voltage with different incident angle $\alpha$.

Figure 3. (a) $P_{v}$ as a function of $V$ for different numbers of superlattice periods $N$. As indicated in figure 3(a) $N$ changes from 2 to 14 in steps of two. (b) Valley-resolved conductances $G_{++,}$ and $G_{--,}$ (in units of $G_{0} = e^{2}N (E)/h$) with $N = 10$ versus $V$. Other parameters are the same as in figure 2.
is \( P_v \) oscillating with \( V \). This behavior arises from the phase coherence of the electron wave functions in the transport process. Another interesting characteristic is that the oscillation magnitude of \( P_v \) is greatly enhanced with increasing \( N \). For example, for low \( V \) the oscillation magnitude of \( P_v \) increases from about 90% for \( N = 2 \) to 100% for \( N = 14 \), while for high \( V \), \( P_v \) increases from less than 40% for \( N = 2 \) to about 90% for \( N = 14 \). With a further increase of \( N \), \( P_v \) gradually tends to be saturated. We note that in the range of \( V = 0 \) to 40 meV one can get a fully valley-polarized current, with its polarized direction modulated by \( V \). It has been shown theoretically that in graphene with broken inversion symmetry, the injection of a valley-polarized current will generate a transverse voltage, in a similar way as an inverse spin-Hall effect. Therefore, the valley polarization produced by our proposed device can be detected directly from the Hall measurement in the outgoing region [6]. In fact, in graphene superlattices with broken inversion symmetry, topological currents originating from graphene’s two valleys were predicted to flow in opposite directions [2]. To see the origin of \( P_v \) oscillations clearly, we show the conductance as a function of \( V \) in figure 3(b) in a magnetic silicene structure of \( N = 10 \). \( G_{+,+} \) and \( G_{-,+} \) exhibits oscillatory behavior, somewhat similar to each other but have different phase constants. The conductance is related to the potential energy profiles in the FS, which are different for different valleys. It is the phase difference.

Figure 4. (a) \( P_v \), (b) \( G_{+,+} \), and (c) \( G_{-,+} \) (in units of \( G_0 = e^2N \langle E \rangle/h \)) as a function of \( V \) with \( \Delta_v = \Delta_v \) (solid line), \( \Delta_v = 3\Delta_v \) (dashed line), and \( \Delta_v = 5\Delta_v \) (dotted line). Other parameters are the same as in figure 2.
between oscillatory \( G_{+,+} \) and \( G_{-,+} \) that makes \( P_v \) oscillating with a large magnitude. It is noteworthy that the conductance spectra exhibit an on/off switching effect by tuning the barrier strength. This character is favorable for electrically controllable device applications.

We consider the effect of \( z \Delta \) in the FS on \( P_v \) in figure 4(a). It is clearly seen that, with increasing \( V \), \( P_v \) exhibit oscillatory behavior with its oscillation magnitude greatly enhanced by \( z \Delta \). The peaks (or valleys) have also been shifted towards higher \( V \) and enlarged gradually. These phenomena correspond to the conductance \( G_{++,} \) and \( G_{--,} \), as shown in figures 4(b) and (c). Like figure 3, for the parameters taken here only \( G_{++,} \) and \( G_{--,} \) are finite, resulting in \( P_v = P \), \( G_{++,} \) and \( G_{--,} \) show oscillatory behavior somewhat similar to each other but have different phase constants. As \( \Delta_v \) increases, each resonant band of both \( G_{++,} \) and \( G_{--,} \) narrows and shifts towards higher \( V \), while the nonresonant gap is broadened. The conductances are strongly dependent on the valley. We observe a coexistence of the insulating transmission gap of one valley and the metallic resonant band of the other. Thus, a strong valley polarization can be obtained.

The conductances are also tunable by the incident energy \( E \). For \( E \) higher than the gap \( \Delta = \lambda_{\omega} + \Delta_v \) in the left electrode, unlike figures 3 and 4, \( G_{++,} \) and \( G_{--,} \) can become finite too; thus \( P_v \) may not be equal to \( P_s \). It is useful to know how \( E \) affects \( P_v \) and \( P_s \), which is shown in figure 5(a). \( P_v \) and \( P_s \) show oscillations with increasing \( E \), and \( P_v \) is equal to \( P_s \) with a full polarization until \( E > 15 \) meV. With a further increase of \( E \), \( P_v \) shows a damped oscillation behavior. At high energy \( P \) practically disappears, while \( P_s \) becomes a small finite value. To understand the behaviors of \( P_v \) and \( P_s \), we study \( G_{\eta,\sigma} \) in figure 5(b). For low energy \( G_{\eta,\sigma} \) does not increase with \( E \), while for high energy \( G_{\eta,\sigma} \) exhibits a nearly linear increase behavior with small oscillation. We can understand these features from \( E_{\eta,\sigma} \) given by equation (2).

\[
|\eta \sigma \lambda_{\omega} - \Delta_v| + V - \alpha h < E < |\eta \sigma \lambda_{\omega} - \Delta_v| + V - \alpha h G_{\eta,\sigma}
\]

originates from the evanescent waves, and electrons can tunnel resonantly through the device at some energies, while \( E > |\eta \sigma \lambda_{\omega} - \Delta_v| + V - \alpha h G_{\eta,\sigma} \) comes from the contribution of the propagating waves, which increases with \( E \) due to the increase of DOS \( N(E) \) in the electrodes.

We now turn to investigate the effect of periods \( N \) of the superlattice on the spin- and valley-dependent thermopowers. It is easily seen from figure 6(a) that \( S_\eta \) is an odd function of \( V \) and presents an oscillation behavior with its magnitude and sign modulated by \( V \). The oscillation magnitude of \( S_\eta \) is greatly enhanced with increasing \( N \) and gradually tends to be saturated at large \( N \). Unlike the charge thermopower, \( S_\eta \) (figure 6(b)) or \( S_\eta \)
Figure 6. (a) $S_c$, (b) $S_s$, and (c) $S_v$ as a function of $V$ for different numbers of superlattice periods $N$. $\Delta_s = 0$ meV in the NS and $\Delta_s = 8.9$ meV in the FS, $D = L = 100$ nm, $E_0 = 0$ meV, $T = 10$ K, $\lambda_{st} = 5.9$ meV, and $h = 5$ meV.

(figure 6(c)) is an even function of $V$. $S_c$ ($S_s$) is also an oscillatory function of $V$. With increasing $N$, $S_c$ ($S_s$) at zero $V$ exhibits a conversion from a dip to a peak. $S_c$ ($S_s$) increases with $N$ and then begins to saturate at large $N$. At zero $V$, a pure $S_c$ ($S_s$) but no $S_v$ can be generated. This is very different from the results in reference [33], where $S_v$ is so weak that it may be overwhelmed by the accompanied $S_c$ of several orders larger. We can explain these features from $S_{\eta \sigma}$ plotted in figures 7(a) and (b). From the symmetry of $E_{s,+}(V) = -E_{s,-}(-V)$ and $E_{s,-}(V) = -E_{s,+}(-V)$, we have $S_{s,+}(V) = -S_{s,-}(-V)$ (figure 7(a)) and $S_{s,-}(V) = -S_{s,+}(-V)$.
(figure 7(b)), which leads to \( S_{+(v)}(V) = S_{-(v)}(-V) \) and \( S_{-}(V) = -S_{+}(-V) \). \( S_{\eta,\sigma} \) exhibits an oscillatory function of \( V \) and can be enhanced by \( N \).

\( S_{c}, S_{s}, \) and \( S_{v} \) are also plotted as a function of \( V \) with different \( z \Delta \) in the FS. As shown in figure 8, with increasing \( z \Delta \), \( S_{c}, S_{s}, \) and \( S_{v} \) exhibit a richer structure. Due to the symmetry of the system, \( S_{c} \) is an odd function of \( V \), while \( S_{s} \) and \( S_{v} \) are an even function of \( V \). As we can see from the inset of figure 8(a), at zero \( V \), \( S_{s} \) and \( S_{v} \) increase nonmonotonically with \( z \Delta \), but \( S_{c} \) is always zero regardless of \( z \Delta \). With the increase of \( V S_{c}, S_{s}, \) and \( S_{v} \) reveal oscillatory behaviors. For fixed \( V \), \( S_{s} \) and \( S_{v} \) nonmonotonically change with \( z \Delta \), with their oscillation magnitudes increasing or decreasing with \( z \Delta \) (see figures 8(a) and (b)). Unlike \( S_{c} \) and \( S_{s} \) in the range of \( z \Delta \), we consider here that \( S_{v} \) reveals a nearly monotonic increase with \( z \Delta \) (figure 8(c)). Those can be understood from \( S \), \( \eta_{\sigma} \). As has been discussed above in figure 7, \( S_{\eta,\sigma} \) oscillates with \( V \) and has different oscillation magnitudes and phase constants for different spins and valleys. \( S_{c}, S_{s}, \) and \( S_{v} \) originate from the contribution of \( S_{++}, S_{+-}, S_{-+}, \) and \( S_{--} \), so due to the combined effect of \( V \) and \( z \Delta \) on \( S_{\eta,\sigma} \), the behaviors of \( S_{c}, S_{s}, \) and \( S_{v} \) become complicated.

Finally, \( S_{c}, S_{s}, \) and \( S_{v} \) are plotted as a function of the Fermi energy \( E_{F} \). As shown in figure 9, \( S_{c} \) is an odd function of \( E_{F} \) and becomes zero at \( E_{F} = 0 \). When \( E_{F} \) deviates from zero, \( S_{c} \) varies rather sharply, changes sign in the symmetry point \( E_{F} = 0 \), and then reaches the maxima on one side of \( E_{F} = 0 \) and the minima on the other side. When \( E_{F} \) is far away from zero, \( |S_{c}| \) declines with \( E_{F} \) and becomes zero at high \( E_{F} \). However, in contrast to \( S_{c}, S_{s}, \) and \( S_{v} \), \( S_{s} \) and \( S_{v} \) are an even function of \( E_{F} \). A maximum magnitude is reached at (or close to) zero \( E_{F} \). By increasing \( E_{F} \), both \( S_{s} \) and \( S_{v} \) can decrease to zero with a small oscillation magnitude. In this case in order to have high thermopowers, we should take a low \( E_{F} \) value.

Although valley-resolved thermoelectric transport through periodic magnetic silicene superlattices gives only a qualitative picture in this work, enhanced valley polarization and thermopowers should be presented in the more realistic cases. For superlattices, the interface roughness can cause some changes of the band structure of the corresponding profile, so they are related to the transmission through the structure. However, differently
from simple NS/FS junctions, the results in periodic superlattices are mainly determined by the periodicity of the barrier potential. This is because the periodic potential leads to resonant bands and nonresonant gaps in the transmission spectra. Therefore, enhanced valley polarization and thermopowers can be achieved that will be insensitive to the interface roughness. This is consistent with the results reported in reference [34], where the authors demonstrated that a moderate disorder would not seriously deteriorate transport and spin-polarization properties of the device.

**Figure 8.** (a) $S_v$, (b) $S_s$, and (c) $S_c$ as a function of $V$ under $D = L = 50$ nm with $\Delta_v = \Delta_w$ (solid line), $\Delta_v = 2\Delta_w$ (dashed line), $\Delta_v = 3\Delta_w$ (dotted line), and $\Delta_v = 4\Delta_w$ (dashed-dotted line). The inset of figure 8(a) shows $S_v$ (solid line), $S_s$ (dashed line), and $S_c$ (dotted line) under $V = 0$ meV versus $\Delta_v$. Other parameters are the same as in figure 6.
4. Summary

In this work we study the valley-resolved thermoelectric transport through a 2D magnetic silicene superlattice. Due to the coupling of spin and valley DOF, the valley- and spin-resolved resonant bands are formed in the transmission spectra, which are separated by nonresonant gaps. The position of each resonant band in the $K$ valley is quite different from that for electrons in the $K'$ valley, which results in strong $P_v$. $P_v$ oscillates with $V$, controlled by the gate voltage, and its magnitude is greatly enhanced by the superlattice structure. In addition, a full valley polarization and an on/off switching effect in the conductance spectra are observed. The effect of $\Delta_0$ in the FS and $E$ on $P_v$ is also considered. Furthermore, the spin- and valley-dependent thermopowers can be controlled by $V$ and Fermi energy, and their magnitudes are greatly enhanced by the superlattice structure. The superlattice structure leads to an enhanced valley-resolved thermoelectric transport with its magnitude controlled by means of gate voltages, making the magnetic silicene superlattice ideal for future valleytronics applications.

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Figure 9. The dependence of $S_c$, $S_s$, and $S_v$ on $E_F$ with $N = 9$. Other parameters are the same as in figure 6.
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