Single- or multi-flavor Kondo effect in graphene

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Abstract – Based on the tight-binding formalism, we investigate the Anderson and the Kondo model for an adatom magnetic impurity above graphene. Different impurity positions are analyzed. Employing a partial-wave representation we study the nature of the coupling between the impurity and the conducting electrons. The components from the two Dirac points are mixed while interacting with the impurity. Two configurations are considered explicitly: the adatom is above one atom (ADA), the other case is the adatom above the center of the honeycomb (ADC). For ADA the impurity is coupled with one flavor for both \(A\) and \(B\) sublattice and both Dirac points. For ADC the impurity couples with multi-flavor states for a spinor state of the impurity. We show, explicitly for a 3\(d\) magnetic atom, \(d_{z^2}\), \((d_{x^2-y^2},d_{xy})\) couple respectively with the \(\Gamma_1, \Gamma_5(E_1)\), and \(\Gamma_6(E_2)\) representations (reps) of \(C_{6v}\) group in ADC case. The bases for these reps of graphene are also derived explicitly. For ADA we calculate the Kondo temperature.

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Introduction. – In graphene, a monolayer of carbon atoms recently fabricated successfully [1], the valence and the conduction bands touch at two inequivalent Dirac points \(K_-\) and \(K_+\) at the corners of the first Brillouin zone (FBZ). Near \(K_-\) and \(K_+\) the low-energy dispersion is linear, indicating a massless Dirac fermions behavior. This particular band structure is at the heart of a number of unusual electronic properties [2]. Graphene is also an interesting candidate for transport applications, in particular for spintronics: It exhibits a remarkably high mobility and the carrier density is controllable by a gate voltage; the mean free path can be as large as 1 \(\mu\)m. Graphene is however not immune to disorder that influences its electronic properties [2]. Extrinsic disorder is realized in a variety of ways: adatoms, vacancies, charges on top of graphene or in the substrates, and extended defects such as cracks and edges. When localized magnetic impurities are added [3] the Kondo effect, i.e. the dynamic screening of the localized moment, has to be addressed at temperatures \(T\) below the Kondo temperature \(T_K\). In this context, previous studies addressed the influence of magnetic impurities using the Hatree-Fock approximation [4,5] which is valid at \(T > T_K\) (see also [6]). In ref. [7] the anisotropic single-channel Kondo model is investigated briefly and in ref. [8] the infinite-\(U\) Anderson model for an impurity embedded in a graphene sheet has been employed and concluded a Fermi liquid behaviour. Recently, it has been claimed that a two-channel Kondo indication in graphene is present due to the valley degeneracy of the Dirac electrons [9], which leads to an over-screening and thus to a non-Fermi-liquid–like ground state. Very recently refs. [10,11] reported on scanning tunneling spectroscopy studies to investigate the Kondo effect in graphene. All of these studies consider a particular configuration of the impurity. In this work, we show that the position of the impurity on or in graphene plays a subtle role and affects essentially the underlying physics. In addition to the role of the vanishing density of state (DOS) at the Dirac points and the linear spectrum in their vicinity, a further important issue is to clarify whether a single or a multi-channel Kondo problem is realized. This is insofar important as the ground states for these two cases are essentially different: As established [12], when the channel number \(N_{cha}\) is equal to \(2S_{im}\) (\(S_{im}\) is the spin of the impurity), the impurity spin is then compensated by the host electrons completely, resulting in a singlet Fermi-liquid ground state. When \(N_{cha} > 2S_{im}\)

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we enter the over-compensated regime where an opposite spin to the original spin appears and acts as a remainder spin; an effective antiferromagnetic coupling between this remainder effective impurity spin and the conducting spins results in a non-Fermi-liquid ground state [13–15]. Thus, it is crucial to clarify when graphene with an impurity with a spin one-half is a single or a two-channel Kondo system. Recently, scanning tunneling conductance spectra are investigated with respect to the position of the scanning tip in refs. [6,16] (see also ref. [17]).

As a first essential step we identify the physical origin of the multi-channels for the Kondo effect by analyzing the tight-binding Anderson model [18,19]. A partial-wave method is then employed to reduce the redundant flavors and to identify the coupling flavors. The role of the Dirac points is exposed. An adatom above one carbon atom (ADA) or an adatom above the center of the honeycomb (ADC) configurations are considered.

**Theoretical formulation.** We start from the Anderson Hamiltonian

\[
H = H_g + H_f + H_{hyb},
\]

(1)

\(H_g (H_f)\) describes graphene (impurity) and \(H_{hyb}\) stands for the impurity-graphene hybridization. The graphene primitive vectors are \(a_1 = a (\frac{1}{2}, \frac{\sqrt{3}}{2})\), \(a_2 = a (\frac{1}{2}, -\frac{\sqrt{3}}{2})\), where \(a = \sqrt{3}a_{cc}\), \(a_{cc}\) is the distance of the nearest carbon atoms. The B sublattice is related to the A sublattice by \(\delta_1 = (a_1 - a_2)/3\), \(\delta_2 = a_1/3 + 2a_2/3\), and \(\delta_3 = -\delta_1 - \delta_2 = -2a_1/3 - a_2/3\) (cf. fig. 1(a)). The two inequivalent Dirac points in the FBZ are \(K_{\pm} = \pm \frac{2\pi}{a}(2/3, 0)\). The second quantization tight-binding-Hamiltonian [2,4,20–22] reads

\[H_g = \sum_{\langle ij \rangle \sigma} (t_{ij} a_{i\sigma} b_{j\sigma} + c^{\dagger}_{i\sigma} V_{i\sigma} c_{j\sigma}),\]

where \(\langle ij \rangle\) is the nearest-neighbor pairs \((i,j)\), \(\sigma\) is a spin index, and \(a_{i\sigma}, b_{i\sigma}\) are annihilation operators for states on the A (B) sublattice. The Hamiltonian in momentum space is

\[H_g = \sum_{\mathbf{k}\sigma} [\xi(\mathbf{k}) a_{\mathbf{k}\sigma}^{\dagger} b_{\mathbf{k}\sigma} + \xi^*(\mathbf{k}) b_{\mathbf{k}\sigma}^{\dagger} a_{\mathbf{k}\sigma}],\]

where \(\xi(\mathbf{k}) = -t\Phi(\mathbf{k})\), \(\Phi(\mathbf{k}) = \sum_\delta e^{i\mathbf{k}\cdot\delta}\), and \(\mathbf{k}\) are wave vectors in the FBZ. The hybridization Hamiltonian in the tight-binding formalism is

\[H_{hyb} = \sum_{\mathbf{k}\sigma} [V_{1\mathbf{k}\sigma} a^{\dagger}_{\mathbf{k}\sigma} f_{\sigma} + V_{2\mathbf{k}\sigma} b^{\dagger}_{\mathbf{k}\sigma} f_{\sigma} + h.c.],\]

where

\[V_{1\mathbf{k}} = \frac{1}{2V} \sum_{\mathbf{R}_s \neq \mathbf{R}_m} e^{-i\mathbf{k} (\mathbf{R}_s + \tau_\alpha) - \mathbf{R}_m} V^{\alpha\mathbf{f}} (\mathbf{R}_s + \tau_\alpha - \mathbf{R}_m),\]

\(\alpha = A\) or \(B\), \(\tau_\alpha\) stands for the relative location for \(\alpha\) atom in the unit cell,

\[V^{\alpha\mathbf{f}} (\mathbf{R}_s + \tau_\alpha - \mathbf{R}_m) = \int d\mathbf{r} [\phi^\alpha (\mathbf{r} - \mathbf{R}_s - \tau_\alpha)]^* h(\mathbf{r}) \times \phi_\mathbf{f} (\mathbf{r} - \mathbf{R}_m),\]

and \(\phi^\alpha\) is the atomic function for the atoms in \(\alpha\) sublattice (only \(\pi\) orbitals are included), \(\phi_\mathbf{f}\) is the localized impurity wave function, \(h\) is the single-particle Hamiltonian, and \(N\) is the number of unit cells.

\[V^{\alpha\mathbf{f}} (\mathbf{R}_s + \tau_\alpha - \mathbf{R}_m)\] is a Slater-type bond in the LCAO formalism [23] which describes the strength of hybridization between two atomic orbitals located at the impurity and its neighbors. Here only \(p_{\pi}\) electrons are relevant. They form the \(\pi\) and \(\pi^*\) which touch at the Dirac points \(K_{\pm}\). An \(s\)-wave impurity substituting one carbon atom in the graphene plane is decoupled from \(\pi\) and \(\pi^*\) bands; for \(d\) - and \(f\)-wave impurity in this substitution case is possible. Our focus is however on the case where the impurity is above the graphene plane. The relative positions of the impurity determine the hybridization strength, \(i.e.\) \(V^{\alpha\mathbf{f}}\) and the phase is included in the exponential in \(V^{\alpha\mathbf{f}}\).

To obtain an effective low-energy Hamiltonian we expand the wave vector \(\mathbf{k} = K_{\pm} + \mathbf{k}\) around \(K_{\pm}\). Thus, \(\Phi(\mathbf{k}) \approx (-v_{\mathbf{F}}/t)(\pm k_x - ik_y)\), or \(\Phi(\mathbf{k})\) is \((v_{\mathbf{F}}/t)k \lambda e^{-i\lambda \theta}\), where \(\lambda = \text{sgn}(s) = 1\) (for \(s = K_+ (K_-)\), \(v_{\mathbf{F}}\) is graphene Fermi’s velocity, \(\theta\) is the azimuthal angle of \(\mathbf{k}\). The Hamiltonian of graphene is then

\[H_g = \sum_{\mathbf{s}, \mathbf{k}} v_{\mathbf{F}} k (a_{\mathbf{s}\mathbf{k}\sigma}^{\dagger} b_{\mathbf{s}\mathbf{k}\sigma}^{\dagger}) \left( \begin{array}{cc} 0 & \lambda e^{-i\lambda \theta} \\ \lambda e^{i\lambda \theta} & 0 \end{array} \right) \left( \begin{array}{c} a_{\mathbf{s}\mathbf{k}\sigma} \\ b_{\mathbf{s}\mathbf{k}\sigma} \end{array} \right),\]

(2)

\(H_g\) is diagonalized by introducing the \(\zeta\) fields [24] as

\[a_{\mathbf{s}\mathbf{k}\sigma} = \frac{1}{\sqrt{2}} [\zeta_{+\mathbf{\sigma}}(\mathbf{k}) + \zeta_{-\mathbf{\sigma}}(\mathbf{k})],\]

\[b_{\mathbf{s}\mathbf{k}\sigma} = \frac{\lambda}{\sqrt{2}} e^{i\lambda \theta} [\zeta_{+\mathbf{\sigma}}(\mathbf{k}) - \zeta_{-\mathbf{\sigma}}(\mathbf{k})],\]

(3)

\[H_g = \sum_{\mathbf{s}, \mathbf{k}} \left[ v_{\mathbf{F}} k |\zeta_{+\mathbf{\sigma}}(\mathbf{k})\zeta_{+\mathbf{\sigma}}(\mathbf{k}) - v_{\mathbf{F}} k |\zeta_{-\mathbf{\sigma}}(\mathbf{k})\zeta_{-\mathbf{\sigma}}(\mathbf{k}) \right].\]

**Partial-wave method and hybridization Hamiltonian.**

Going over from a discrete to a continuum \(\mathbf{k}\) [25] we write

\[H_g = \int d\mathbf{k} \left[ v_{\mathbf{F}} k |\zeta_{+\mathbf{\sigma}}(\mathbf{k})\zeta_{+\mathbf{\sigma}}(\mathbf{k}) - v_{\mathbf{F}} k |\zeta_{-\mathbf{\sigma}}(\mathbf{k})\zeta_{-\mathbf{\sigma}}(\mathbf{k}) \right],\]

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and
\[
H_{\text{hyb}} = \frac{(N\Omega_0)^{1/2}}{2\pi} \sum_{s\sigma} \int \frac{dk}{2\pi} \left[ (V_{sk}^A a_{sk\sigma}^\dagger f_{\sigma} + V_{sk}^B b_{sk\sigma}^\dagger f_{\sigma}) + h.c. \right].
\]

\(\Omega_0\) is the unit cell area. In 2D orbital momentum eigenfunctions \(e^{im\theta}\) we write
\[
\zeta^{\pm}_{s\sigma}(k) = \frac{1}{\sqrt{|k|}} \int_{m=-\infty}^{\infty} \frac{1}{\sqrt{2\pi}} e^{im\theta} \zeta^{ms}_{s\sigma}(k),
\]
where
\[
\{\zeta^{ms}_{s\sigma}(k), \zeta^{ms\dagger}_{s\sigma}(k')\} = \delta_{s\sigma} \delta_{ss'} \delta_{mm'} \delta(k-k').
\]

As we set the spatial zero point "O" as the projection of impurity position onto the unit cell. 

\textbf{Adatom impurity above one A atom (ADA):} In this case \(\tau_A = 0\), in the Dirac cones we find (we assume \(v_0\) is real) \(V_{sk}^A = V_{sk}^A(0) = \frac{\lambda}{\sqrt{2\lambda}}\). The nearest neighbors are three B atoms at \(\delta_i, i = 1, 2, 3\). Thus, \(V_{sk}^B = \Phi^s(k)\), where \(v_1 = V_{sk}^B = \frac{\lambda}{\sqrt{2\lambda}}\). Substituting in \(H_{\text{hyb}}\), we find
\[
H_{\text{hyb}} = \int_{m=-\infty}^{\infty} v_F \left[ \left( \zeta^{ms\dagger}_{s\sigma}(k) \zeta^{ms}_{s\sigma}(k) - \zeta^{ms\dagger}_{s\sigma}(k) \zeta^{ms}_{s\sigma}(k) \right) \right] f + h.c.
\]
This equation tells us only \(m = 0\) flavor couples to the impurity.

\textbf{Adatom impurity above the center of the honeycomb:} Here, the impurities hybridize with the same strength \(v_0\), with 6 nearest neighbors, 3 A atoms and 3 B atoms. With respect to O, A2, B2, A3, B3 are in other unit cells. Three \(A\) (\(B\)) atoms yield \(\Phi^s(\Phi)\) when the phases are coherently superimposed, i.e.
\[
V_{sk}^A a_{sk\sigma}^\dagger f_{\sigma} + V_{sk}^B b_{sk\sigma}^\dagger f_{\sigma} = \frac{v_0}{\sqrt{2\lambda}} \left[ e^{-ik\delta_1} (1 + e^{-ik(a_1+a_2)}) + e^{-ik(a_1)} a_{sk\sigma}^\dagger f_{\sigma} + e^{-i(k\delta_1)} (1 + e^{-ik(a_1+a_2)}) b_{sk\sigma}^\dagger f_{\sigma} \right]
\]
For \(H_{\text{hyb}}\) we deduce
\[
H_{\text{hyb}} = \int_{m=-\infty}^{\infty} v_F \left[ \left( \zeta^{ms\dagger}_{s\sigma}(k) \zeta^{ms}_{s\sigma}(k) - \zeta^{ms\dagger}_{s\sigma}(k) \zeta^{ms}_{s\sigma}(k) \right) \right] f + h.c.
\]
When \(s = K_+, \lambda = 1\), the impurity couples with the partial wave \(m = 1\) from \(A\) sublattice, and \(m = -2\) from \(B\) sublattice. When \(s = K_-\), \(\lambda = -1\), it interacts with \(m = -1\) from \(A\) sublattice, and \(m = 2\) from \(B\) sublattice.

\textbf{Flavor right movers and symmetry analysis.} We unfold the range of momenta \(k\) from \((0, \infty)\) to \((-\infty, +\infty)\) by defining flavor right movers. For the ADA case they are \(c^{s\sigma}_1(k) = \zeta^{s\sigma}_1(k)\), for \(k > 0\) and \(c^{s\sigma}_0(k) = \zeta^{s\sigma}_0(k)\), for \(k < 0\). We obtain thus
\[
H_{\text{hyb}} = \int_{m=-\infty}^{\infty} v_F \left| k \right| f + h.c.
\]
where \(\epsilon_k = h v_F k\). So, the impurity couples with only one flavor of conducting waves. For the ADC case the impurity couples with \(m = \lambda\) or \(m = -2\lambda\) flavors from \(A\) or \(B\) sublattices, respectively. Therefore, we introduce two flavor right movers as
\[
c^s_1 = \lambda \zeta^{s\sigma}_1(k), \quad c^s_2 = \zeta^{s\sigma}_2(k)
\]
for \(k > 0\);
\[
c^s_1 = -\lambda \zeta^{s\sigma}_1(k), \quad c^s_2 = \zeta^{s\sigma}_2(k)
\]
for \(k < 0\).

In terms of these, the Hamiltonian is
\[
H_{\text{hyb}} = \int_{m=-\infty}^{\infty} v_F \left| k \right| f + h.c.
\]
and
\[
H_{\text{hyb}} = \int_{m=-\infty}^{\infty} v_F \left| k \right| f + h.c.
\]
where \(n = 1, 2\). In these two cases, we have \(c^{s\sigma}_n(k), c^{s\dagger\sigma}_n(k')\) are the impurity couples with the partial wave \(m = 1\) from \(A\) sublattice, and \(m = -2\) from \(B\) sublattice.
Table 1: Character table for double group $C_{3v}$ for the $K_\pm$ point of graphene. $\Gamma_{1,5,6}$ are extra representations (reps) and the spinor rep is $\Gamma_4$.

| $C_{3v}$ | $E$ | $E'$ | $C_3^+$, $C_3^-$ | $C_3^{+*}$, $C_3^{-*}$ | $\sigma_{d_1}$ | $\sigma_{d_2}$ |
|----------|-----|------|-----------------|-------------------|----------------|----------------|
| $\Gamma_1$ | 1 | 1 | 1 | 1 | 1 |
| $\Gamma_2$ | 1 | 1 | 1 | 1 | -1 |
| $\Gamma_3$ | 2 | 2 | -1 | -1 | 0 |
| $\Gamma_5$ | 1 | -1 | -1 | 1 | -i |
| $\Gamma_6$ | 1 | -1 | -1 | 1 | i |
| $\Gamma_4$ | 2 | -2 | 1 | -1 | 0 |

$E, C_3^+, C_3^-, \sigma_{d_1}$, where $i = 1, 2, 3$. The symmetry operation $\sigma_{d_1}$ leaves the Dirac points unchanged but interchanges $A$ by $B$ atoms in real space, $\sigma_{d_2}$ acts in the opposite way.

If we use $D^\pm$ to indicate the spinor reps for a rotation group, then for $C_{3v}$, $\Gamma_1 \otimes D^\pm = \Gamma_4$. We also have the relations $\Gamma_5 \otimes \Gamma_4 = \Gamma_3$, $\Gamma_6 \otimes \Gamma_4 = \Gamma_3$, and $\Gamma_4 \otimes \Gamma_4 = \Gamma_1 + \Gamma_2 + \Gamma_3$. Without spin and when the Fermi level $E_F$ crosses $K_{\pm}$ we have a $\Gamma_4$ reps. Following ref. [28] we can show the Bloch sum for the $p_z$ orbitals on $A$ and $B$ sublattices are the basis of the representation, i.e. $\phi_{\alpha \sigma}^\pm (r) = \frac{1}{\sqrt{N}} \sum \langle \mathbf{r} | \mathbf{R}_i \rangle \phi_{\alpha \sigma}^\pm (\mathbf{r} - \mathbf{R}_i - \mathbf{r}_0)$, where $\alpha = A$ or $B$, $\phi_{\alpha \sigma}^\pm$ representing $p_z$ orbital at $\mathbf{R}_i + \mathbf{r}_0$ site. When spin is taken into account and spin-orbit interaction (SOI) is ignored, the degeneracy at $K_{\pm}$ is $2 \times 2 = 4$. When SOI is considered, the reps of double group are split into $\Gamma_3 \otimes D^{\pm} = \Gamma_4 + \Gamma_5 + \Gamma_6$. The reps $\Gamma_4, \Gamma_5$, and $\Gamma_6$ are degenerate in absence of SOI at $K_{\pm}$.

a) For the ADA case the system has a $C_{3v}$ symmetry (different from $C_{3v}$, the character table is shown by table 1). The planes reflection symmetry in $C_{3v}$ converts the point $K_+ \rightarrow K_-$ and vice versa. Thus, the symmetry $C_{3v}$ mixes the states from the two Dirac points. In fact, this group is a subgroup of $C_{6v}$ and the symmetry reflection planes do show up in $C_{6v}$ and map Dirac points onto each others. Thus, only one flavor couples with the impurity. The Hamiltonian in this case shows that a basis belonging to $\Gamma_1$ representation of $C_{3v}$ can be constructed from the $E$ rep of $C_{3v}$, which leads to an invariant hybridization Hamiltonian. If we consider explicitly a 3d magnetic impurity in this case, for example Mn (as realized in the ADA case [32]), the nonvanishing coupling is only present between the $d_{z^2}$ orbital of the impurity and graphene since the $d_{z^2}$ orbital transforms according to $\Gamma_1$ rep.

b) For the ADC case the symmetry group is $C_{6v}$ where $C_2$, $C_6^0$, $C_6^{\pm*}$ and $\sigma_{d_1}$ show up in this group except the symmetry operations in $C_{3v}$. In this group, $C_2$, $C_6^0$ transform the $A$ sublattice into the $B$ sublattice, $K_+ \rightarrow K_-$ and vice versa. $\sigma_{d_1}$ keep the sublattice unchanged but exchange $K_+$ and $K_-$. Therefore, we have to construct the basis for $C_{6v}$ by using the basis for the two Dirac points in the $C_{3v}$ group. From table 2, $E$ rep in $C_{3v}$ corresponds to the $E_1$ and $E_2$ reps in $C_{6v}$. Therefore we can use the Bloch sums $\phi_{K_\pm \kappa \sigma}^{A(B)}$ in the $E$ rep in $C_{3v}$ to construct the reps in $C_{6v}$. Explicitly we consider a 3d magnetic element, and ignore the crystal electric field splitting so that the energy levels for different reps are still degenerate or quasi-degenerate. The wave function should be constructed according to the symmetry group of $C_{6v}$. The reason is that the nonvanishing hybridization should be invariant under the operations in $C_{6v}$ which is obtained as $\Gamma_1 \otimes \Gamma_1$ by group theory. From table 2, we infer $d_{z^2}$ belongs to $\Gamma_1$, $(d_{x^2-y^2}, d_{xz}, d_{yz})$ to $\Gamma_1(E_1)$, and $(d_{x^2-y^2}, d_{xy})$ to $\Gamma_1(E_2)$. These 3d orbitals are well defined.

Having specified the states of the impurity, we consider the basis of $C_{6v}$ reps by combining the basis of $E$ in $C_{3v}$. We find

$$\phi_{K_\pm \kappa \sigma} = (1/\sqrt{2})((\phi_{K_+ \kappa \sigma}^A - \phi_{K_- \kappa \sigma}^B), -\phi_{K_+ \kappa \sigma}^B + \phi_{K_- \kappa \sigma}^A)^T, \phi_{K_\pm \kappa \sigma} = (1/\sqrt{2})((\phi_{K_+ \kappa \sigma}^A + \phi_{K_- \kappa \sigma}^B), \phi_{K_+ \kappa \sigma}^B + \phi_{K_- \kappa \sigma}^A)^T, \quad (6)$$

where $\omega = e^{i2\pi/3}$. Therefore, the hybridization Hamiltonian for 3d magnetic impurity is

$$H_{hyb} = \sum_{\alpha \sigma \kappa} \phi_{K_\pm \kappa \sigma}^\dagger (\phi_{K_\pm \kappa \sigma}^A + \omega^{*} \phi_{K_\pm \kappa \sigma}^B) + h.c., \quad (7)$$

where $\alpha = \Gamma_1, \Gamma_5, \Gamma_6$ reps, and the latter two are 2D reps. The creation and annihilation operators are obtained by defining the field operators based on the respective basis of the corresponding rep in the standard way. This situation resembles the Cooper-Schrieffer model [33], however the Hamiltonian now is written in terms of the irreducible reps of the systems.

**Kondo temperature under large $U$.** We calculated the Kondo temperature for the Anderson model in the ADA case using the equation of motion method [34] that gives consistent results for the susceptibility with those derived by other methods at high temperature [35] and delivers the correct Kondo temperature [36]. The Kondo temperature for large $U$ is proportional to $\exp(\pi^2 g^2 / 24 \mu)$, where $\Delta = \pi \zeta^2 / 24g^2D^2$ is the energy cutoff. If $\mu < \mu$ and with increasing $|\mu|$ the Kondo temperature increases exponentially. Thus, a small gate voltage may change the Kondo temperature dramatically [9].

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Concluding remarks. — Summarizing, we investigated the Kondo model for a magnetic impurity in graphene in an adatom configuration. Starting from the tight-binding Anderson model we applied a partial-wave method to expose the coupling flavors. The Kondo effect in graphene behaves as follows: i) For the ADA case, only one flavor couples to the impurity. ii) For the ADC case, the number of channels is determined by the multiplets of the impurity as well. Multi-flavor character shows up in this case. iii) Based on the irreducible basis of the reps of the system, a Coqblin-Schrieffer–like hybridization model is derived for a 3d magnetic atom. \(d_{z^2}, (d_{x^2},d_{y^2})\), and \((d_{x^2-y^2},d_{xy})\) couple with the \(\Gamma_1, \Gamma_5(E_1)\), and \(\Gamma_6(E_2)\) reps, respectively. The bases for these reps of graphene are derived explicitly. iv) The degeneracy of the two Dirac points only leads to a higher \(T_K\) for ADA case. Applying a gate voltage the exchange coefficients are non-zero due to a finite DOS at \(E_F\) and the Kondo effect should be observable.

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