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Exchange interactions and frustrated magnetism in single-side hydrogenated and fluorinated graphene

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Magnetism in single-side hydrogenated (C_2H) and fluorinated (C_2F) graphene is analyzed in terms of the Heisenberg model with parameters determined from first principles. We predict a frustrated ground state for both systems, which means the instability of collinear spin structures and sheds light on the absence of a conventional magnetic ordering in defective graphene demonstrated in recent experiments. Moreover, our findings suggest a highly correlated magnetic behavior at low temperatures offering the possibility of a spin-liquid state.

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Magnetism of sp materials is an exotic phenomenon of a fundamental physical relevance and potential practical importance. In particular, one can expect much higher Curie temperatures for sp magnets than for conventional magnetic semiconductors [1]. A discovery of graphene has revived an interest in this subject [2, 3]. Recent studies provided persuasive experimental evidence for the existence of magnetism in graphene doped by hydrogen and fluorine atoms [4–6]. Although no magnetic ordering has been detected in defective graphene even at liquid helium temperatures [6, 7], theoretical studies show that an ordered decoration of graphene with a composition of C:H(F)=2:1 (occasionally referred to as graphone) is unstable toward the formation of (anti-)ferromagnetism [8, 9], offering perspectives of applications [10].

In contrast to single defects in graphene, whose magnetic properties are theoretically well understood [11, 12], the magnetism in C_2H and C_2F lacks a consistent description. Although previous density functional (DF) studies point to the possibility of magnetic ordering in C_2H and C_2F [8–10], its origin and stability have not yet been analyzed in details. Moreover, the prediction of the magnetic ground state cannot be conclusive within the DF scheme due to symmetry constraints imposed by periodic boundary conditions. On the other hand, the possibility of a description in terms of spin models is not obvious due to the absence of a unique definition of local magnetic moments in sp systems, which makes the parametrization of a model particularly challenging.

In this Letter, we construct an appropriate Heisenberg model to describe magnetic properties of C_2H and C_2F , and determine its parameters from first principles. For both systems, we find a frustrated ground state, which contradicts previous theoretical predictions. We show that the possibility of spontaneous magnetization in C_2H and C_2F is very unlikely, instead, an unconventional magnetic behavior is expected at low temperatures, such as

non-collinear ordering or a spin liquid state.

We first analyze the electronic structure at the DF level and construct a tight-binding Hamiltonian by taking advantage of the Maximally-Localized Wannier Functions (MLWF) [13]. DF calculations were performed within the plane-wave pseudopotential method as implemented in the QUANTUM-ESPRESSO simulation package [14]. Exchange-correlation effects were taken into account by using the standard GGA-PBE functional [15]. To achieve high accuracy, we employed an energy cutoff of 50 Ry for the plane-wave basis and 400 Ry for the charge density, as well as a (64×64) k -point mesh. The surface layers were separated by a vacuum region of

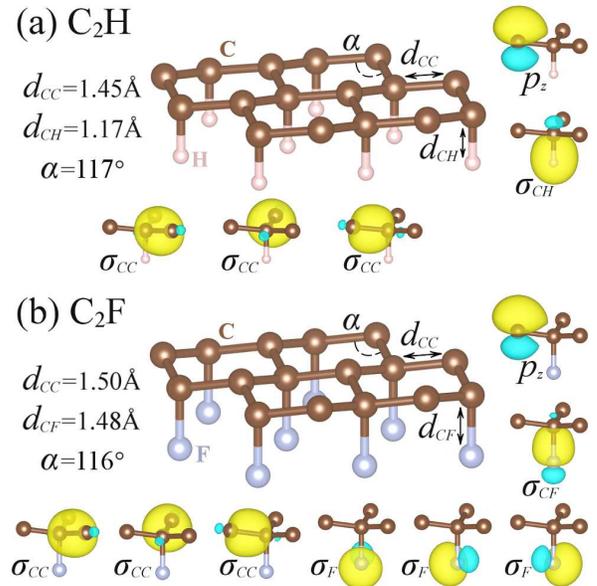


FIG. 1. (Color online) Schematic representation of the atomic structure and Maximally-Localized Wannier Functions for C_2H and C_2F .

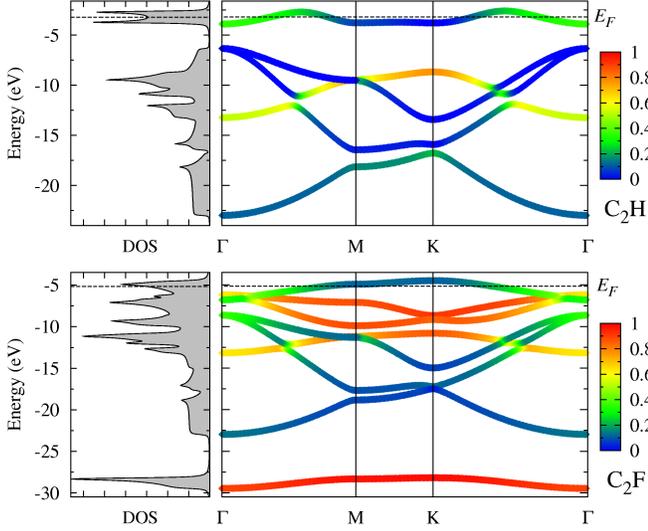


FIG. 2. (Color online) Band structure and density of states (DOS) for C_2H and C_2F obtained from non-spin-polarized calculations. Contributions from the impurity orbitals (s for H and sp for F) are shown by color.

40 Å and fully relaxed. To construct a localized basis for the electronic Hamiltonian, we used the WANNIER90 code [16].

Fully relaxed structures of C_2H and C_2F as well the MLWF are shown in Fig. 1. All orbitals are well localized providing an intuitive picture of the chemical bonding. In each case, the MLWF basis includes three bond-centered orbitals, corresponding to the sp^2 -like hybridization of carbon atoms (σ_{CC}), one bond-centered sp -like orbital localized between carbon and hydrogen/fluorine atoms (σ_{CH}/σ_{CF}), and one p_z -like orbital centered at the non-bonded carbon atom. In the case of C_2F , there are three additional orbitals resulting from the F-centered sp^2 -type hybrids (σ_F). In contrast to C_2H , the σ_{CF} and p_z orbitals in C_2F are more localized, implying stronger binding and enhanced stability of the fluorinated system (Table I).

In Fig. 2, we show the density of states (DOS) and band structure projected onto the orbitals associated with the MLWF localized at the impurity orbitals of C_2H and C_2F . In both cases there is a density of states at the Fermi level, indicating the presence of uncompensated electrons, which are localized mainly at the p_z -orbital of carbon with some contribution from the impurity orbitals.

The MLWF formalism allows us to represent the electronic Hamiltonian in the tight-binding form,

$$H = \sum_i \varepsilon_i n_i + \sum_{i \neq j} t_{ij} c_i^\dagger c_j, \quad (1)$$

where summation runs over the Wannier orbitals, ε_i is the energy of the i th orbital, t_{ij} is the hopping parameter between i th and j th MLWFs, and c_i^\dagger (c_j) is the creation

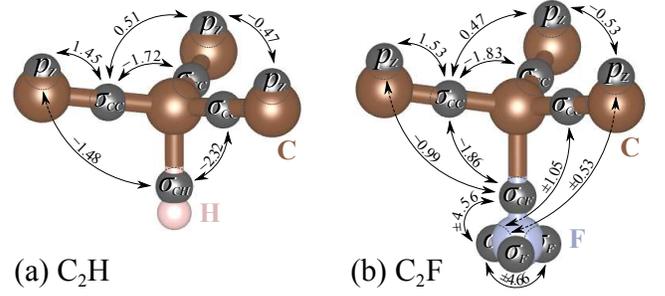


FIG. 3. (Color online) The most relevant hopping parameters (in eV) between different Wannier orbitals for C_2H (a) and C_2F (b). The orbitals are shown schematically by gray labelled spheres, localized at the corresponding MLWF centers. Positive and negative values correspond to destructive and constructive overlap between the involved orbitals, respectively.

(annihilation) operator of electrons localized at the i th (j th) MLWF. In Fig. 3, we show the most relevant hopping parameters obtained from the non-spin-polarized Hamiltonian for C_2H and C_2F . In contrast to graphene, where nearest-neighbor hoppings only are essential to describe the main characteristics of the ground-state electronic spectrum, the case of graphene derivatives is considerably more complicated due to the lowering of symmetry. In view of the complexity of the Hamiltonian, a simple analytical model, which would be appropriate for studying magnetic properties of C_2H and C_2F , is not readily available.

If the spin polarization is switched on, the exchange interactions split the energies of electrons with opposite spins, resulting in the formation of magnetic states. Interestingly, spin-polarized DF calculations yield different ground states for C_2H and C_2F , corresponding to the ferromagnetic (FM) and antiferromagnetic (AFM) configurations with a small band gap (~ 0.5 – 1.0 eV), respectively [8–10]. Exploiting exponential localization of the MLWF in real space [17], the magnetic moments can be associated with the respective Wannier orbitals. Accordingly, the exchange splitting on the i th Wannier orbital can be calculated from diagonal elements of the spin-polarized MLWF Hamiltonian as $\Delta_i = H_{ii}^\uparrow - H_{ii}^\downarrow$, whereas the corresponding magnetic moment is easily obtained from the density of states (DOS) $g_i^\sigma(\varepsilon)$ projected onto the i th MLWF, $M_i = \int_{-\infty}^{E_F} d\varepsilon [g_i^\uparrow(\varepsilon) - g_i^\downarrow(\varepsilon)]$. The corresponding values are listed in Table I for different Wannier orbitals. Apart from the magnetic moment localized in the p_z orbital, there is a non-negligible contribution from the impurity orbitals, which is more pronounced for C_2H (σ_{CH} orbital). Below, we show that this contribution plays an essential role in determining the magnetic behavior of C_2H .

Owing to the possibility of a localized representation of magnetic moments in C_2H and C_2F , the magnetic in-

TABLE I. Linear spreads (Ω), magnetic moments (M) and exchange splittings (Δ) associated with the Wannier orbitals of C_2H and C_2F . The orbitals are denoted according to Fig. 1. The spreads are given for the non-spin-polarized orbitals.

	C_2H			C_2F			
	σ_{CC}	σ_{CH}	p_z	σ_{CC}	σ_{CF}	p_z	σ_F
Ω , Å	0.86	1.02	1.25	0.84	0.79	1.11	0.71
$ \Delta $, eV	0.29	0.62	2.54	0.24	0.41	1.92	0.24
M , μ_B	0.02	0.17	0.76	0.01	0.05	0.59	0.02

interactions can be analyzed in terms of the Heisenberg model,

$$H = - \sum_{i \neq j} J_{ij} \mathbf{s}_i \cdot \mathbf{s}_j, \quad (2)$$

where J_{ij} is the exchange coupling between i and j sites, and \mathbf{s}_i , \mathbf{s}_j are unit vectors pointing in the direction of local magnetic moments at sites i and j , respectively. Thus, we consider the mapping of DF calculations onto *classical* Heisenberg model [18]. At this mapping, the exchange parameters can be calculated via the second variations of the total energy with respect to infinitesimal rotations of the magnetic moments. For a lattice with a basis, this can be written in the following form [18]

$$J_{ij}^{\alpha\beta} = \frac{1}{4\pi} \int_{-\infty}^{E_F} d\varepsilon \text{Im} \left[\Delta_{\alpha} G_{ij}^{\alpha\beta\downarrow}(\varepsilon) \Delta_{\beta} G_{ji}^{\beta\alpha\uparrow}(\varepsilon) \right], \quad (3)$$

where i, j and α, β denote unit cell and sublattice (orbital) indices, respectively. $\Delta_{\alpha(\beta)}$ is the exchange splitting of the orbital α (β), E_F is the Fermi level, and $G_{ij}^{\alpha\beta\sigma}(\varepsilon)$ is the real-space Green's function, whose reciprocal-space representation is given by,

$$G_{\mathbf{k}}^{\sigma}(\varepsilon) = [\varepsilon - H^{\sigma}(\mathbf{k}) + i\eta]^{-1}, \quad (4)$$

$\eta \rightarrow +0$. Here, $H^{\sigma}(\mathbf{k})$ is the reciprocal-space Hamiltonian for spin $\sigma = \uparrow, \downarrow$, whose matrix elements are obtained in the MLWF basis from spin-polarized DF calculations, corresponding to the most favorable magnetic state (FM for C_2H , and row-wise AFM for C_2F). The Green's functions are evaluated by taking $\approx 10^6$ \mathbf{k} -points in the full Brillouin zone and the parameter $\eta = 0.02$.

In Fig. 4, we show the calculated exchange parameters for both C_2H and C_2F as functions of the distance between the pairs of the relevant orbitals, i.e., p_z-p_z and $p_z-\sigma_{CH(CF)}$. The other orbitals carry negligible moments and thus do not contribute to the exchange. The leading exchange interaction is different in sign for the two systems. In the case of C_2H , it corresponds to FM interactions between the p_z and $\sigma_{CH(CF)}$ orbitals, whereas for C_2F , it originates from the p_z-p_z AFM coupling. Compared to C_2H , where the leading FM term is followed by a weaker AFM contribution, the interactions in C_2F are almost negligible at larger distances and determined

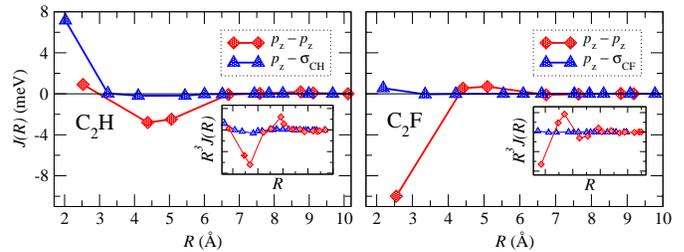


FIG. 4. (Color online) Calculated exchange parameters as functions of the distance between the p_z-p_z (red curve) and $p_z-\sigma_{CH(CF)}$ (blue curve) orbitals in C_2H (left panel) and C_2F (right panel). The interactions between the other orbitals are negligibly small and not shown. In the inset, the $J(R)$ parameters are shown with a prefactor R^3 .

by the first AFM term only. Although for both systems the p_z-p_z exchange parameters change their sign with distance, this interaction term decays faster than R^{-3} (see the inset of Fig. 4). This behavior indicates the absence of RKKY-like interactions, which is not surprising in view of insulating character of both systems. Importantly, while the magnetic moments arise primarily from the p_z orbitals, the dominant contribution to the exchange in C_2H comes from the FM interaction between the spins localized on the unsaturated carbon (p_z) and impurity (σ_{CH}) orbitals. This contribution is therefore responsible for the previously reported difference between the ground states of C_2H and C_2F at the DF level [8–10].

In fact, the difference of an order of magnitude between the nearest-neighbor interactions in C_2H and C_2F can qualitatively be explained by the difference in the overlap between the p_z and $\sigma_{CH(CF)}$ orbitals, resulting from the unequal spreads (Table I) and interorbital distances (Fig. 1). Thus, this kind of exchange interaction is directly related to the overlap of the wave functions involved. The second contribution to the exchange (p_z-p_z interactions) exhibits a more complicated behavior, which is typical for the indirect superexchange interactions. By comparing the hopping parameters given in Fig. 3, it can be deduced that the difference in the p_z-p_z coupling between C_2H and C_2F is related to the presence of the $2p$ fluorine orbitals (σ_F) in C_2F , which provide an additional superexchange pathway compared to C_2H . For both systems, however, a conventional perturbative treatment of the superexchange (e.g., in terms of the relation $J_{ij} \sim t_{ij}^2/U$) is probably not correct due to the small magnitude of the Coulomb interactions (U) in sp systems.

We now turn to the problem of magnetic ordering in C_2H and C_2F . In both cases, the presence of a significant AFM interaction on the triangular lattice of p_z orbitals suggests the instability of any collinear spin alignment [19]. Surprisingly, the exchange interactions in C_2F represent almost an ideal realization of the Heisenberg model with nearest-neighbor AFM couplings ($J = -10$

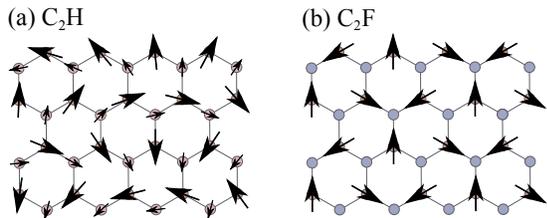


FIG. 5. (Color online) Spin structures corresponding to a classical ground state for: (a) C_2H (incommensurate spin spiral), and (b) C_2F (120° Néel state).

meV), whose classical solution corresponds to the 120° Néel state [Fig. 5(b)]. In contrast to C_2F , the ground state of C_2H is not obvious due to the presence of a larger number of competing interactions of different signs. In this case, we seek the ground state spin configuration in the general form of a coplanar spin spiral, $\mathbf{S}_{\alpha(\beta)}(\mathbf{r}) = S_{\alpha(\beta)} (\cos(\mathbf{Q} \cdot \mathbf{r} + \phi_{\alpha(\beta)}), \sin(\mathbf{Q} \cdot \mathbf{r} + \phi_{\alpha(\beta)}))$, where \mathbf{Q} is the spiral ordering vector, and $\phi = \phi_\alpha - \phi_\beta$ is the phase shift between the spins on different sublattices. Minimizing the total energy with respect to \mathbf{Q} and ϕ we find that the ground state corresponds to an incommensurate spin spiral with $\mathbf{Q}=(0.31,0.79)$ and $\phi=41.5^\circ$. The corresponding spin configuration is shown schematically in Fig. 5(a).

The problem of magnetic frustration in C_2H and C_2F becomes more complicated if the quantum effects are taken into account. If we assume that the unsaturated carbon atoms in C_2F carry $S = 1/2$ spin, its ground state would correspond to the three-sublattice magnetic order, similar to the classical case, which is a widely accepted scenario for the spin-1/2 quantum Heisenberg antiferromagnet on a triangular lattice [20]. However, the magnitude of the local magnetic moment in C_2F ($M_{p_z} \approx 0.59\mu_B$) strongly suggest itinerant-electron magnetic behavior, which cannot be properly described by the quantum Heisenberg model. Although multispin models can provide a consistent description of the magnetism on the triangular lattice [21], the way of the parametrization of such models is far from being clear. The most natural model for studying itinerant-electron magnetism is the Hubbard model. Previous applications of this model to the triangular lattice reveal a spin liquid state at intermediate values of the U/t parameter with the transition to the ordered Néel state at higher U/t values [21, 22]. Given a relatively small magnitude of U in sp materials due to the delocalization of electrons [23], the spin liquid state appears to be a possible scenario for the ground state of C_2F .

The case of C_2H turns out to be even more complicated due to the presence of the two magnetic sublattices comprising the honeycomb lattice. Although recent studies of the Hubbard model on the honeycomb lattice do not exclude the existence of a spin liquid phase in some

range of the U/t parameters [24], their results cannot be directly applied to C_2H having two nonequivalent sublattices with $M_{p_z} \approx 0.76\mu_B$ and $M_{\sigma_{CH}} \approx 0.17\mu_B$. An adequate description of such a system requires a more sophisticated analysis in terms of the multiorbital Hubbard model and represents an important issue for further studies.

In conclusion, we have studied the exchange interactions in C_2H and C_2F and analyzed the possibilities of the magnetic ordering in these systems. In each case, we have found the presence of AFM interactions on the triangular lattice of magnetic moments, leading to the instability of the collinear magnetic ordering due to frustration. Although the true quantum ground state remains to be determined, both systems are expected to exhibit highly correlated magnetic behavior at low temperatures and can be considered as promising candidates for the quantum spin liquid.

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- [1] D. M. Edwards and M. I. Katsnelson, *J. Phys.: Condens. Mat.* **18**, 7209 (2006).
- [2] O. V. Yazyev, *Rep. Prog. Phys.* **73**, 056501 (2010).
- [3] M. I. Katsnelson, *Graphene: Carbon in Two Dimensions* (Cambridge Univ. Press, Cambridge, 2012), Chap. 12.
- [4] X. Hong, K. Zou, B. Wang, S.-H. Cheng, and J. Zhu, *Phys. Rev. Lett.* **108**, 226602 (2012).
- [5] K. M. McCreary, A. G. Swartz, W. Han, J. Fabian, and R. K. Kawakami, *Phys. Rev. Lett.* **109**, 186604 (2012).
- [6] R. R. Nair, M. Sepioni, I-Ling Tsai, O. Lehtinen, J. Keinonen, A. V. Krasheninnikov, T. Thomson, A. K. Geim, and I. V. Grigorieva, *Nat. Phys.* **8**, 199 (2012).
- [7] M. Sepioni, R. R. Nair, S. Rablen, J. Narayanan, F. Tuna, R. Winpenny, A. K. Geim, and I. V. Grigorieva, *Phys. Rev. Lett.* **105**, 207205 (2010).
- [8] J. Zhou, Q. Wang, Q. Sun, X. S. Chen, Y. Kawazoe, and P. Jena, *Nano Lett.* **9**, 3867 (2009).
- [9] D. W. Boukhvalov, *Physica E* **43**, 199 (2010).
- [10] L. Li, R. Qin, H. Li, L. Yu, Q. Liu, G. Luo, Z. Gao, and J. Lu, *ACS Nano* **5**, 2601 (2011).
- [11] O. V. Yazyev and L. Helm, *Phys. Rev. B* **75**, 125408 (2007).
- [12] E. Kogan, *Phys. Rev. B* **84**, 115119 (2011).
- [13] N. Marzari, A. A. Mostofi, J. R. Yates, I. Souza, and D. Vanderbilt, *Rev. Mod. Phys.* **84**, 1419 (2012).
- [14] P. Giannozzi *et al.*, *J. Phys.: Condens. Matter*, **21**, 395502 (2009).
- [15] J. P. Perdew, K. Burke, and M. Ernzerhof, *Phys. Rev. Lett.* **77**, 3865 (1996).
- [16] A. A. Mostofi, J. R. Yates, Y.-S. Lee, I. Souza, D. Vanderbilt, and N. Marzari *Comput. Phys. Commun.* **178**, 685 (2008).
- [17] C. Brouder, G. Panati, M. Calandra, C. Mourougane, and N. Marzari, *Phys. Rev. Lett.* **98**, 046402 (2007)

- [18] A. I. Lichtenstein, M. I. Katsnelson, V. P. Antropov, and V. A. Gubanov, *J. Magn. Magn. Mater.* **67**, 65 (1987).
- [19] G. H. Wannier, *Phys. Rev.* **79**, 357 (1950).
- [20] B. Bernu, P. Lecheminant, C. Lhuillier, and L. Pierre, *Phys. Rev. B* **50**, 10048 (1994); L. Capriotti, A. E. Trumper, and S. Sorella, *Phys. Rev. Lett.* **82**, 3899 (1999); D. J. J. Farnell, R. F. Bishop, and K. A. Gernoth, *Phys. Rev. B* **63**, 220402(R) (2001); S. R. White and A. L. Chernyshev, *Phys. Rev. Lett* **99**, 127004 (2007).
- [21] H.-Y. Yang, A. M. Läuchli, F. Mila, and K. P. Schmidt, *Phys. Rev. Lett.* **105**, 267204 (2010).
- [22] A. Singh, *Phys. Rev. B* **71**, 214406 (2005); P. Sahebsara and D. Senechal, *Phys. Rev. Lett.* **100**, 136402 (2008); A. E. Antipov, A. N. Rubtsov, M. I. Katsnelson, and A. I. Lichtenstein, *Phys. Rev. B* **83**, 115126 (2011).
- [23] M. Schüler, M. Rösner, T. O. Wehling, A. I. Lichtenstein, and M. I. Katsnelson, arXiv:1302.1437.
- [24] Z. Y. Meng, T. C. Lang, S. Wessel, F. F. Assaad, and A. Muramatsu, *Nature (London)* **464**, 847 (2010); S. Sorella, Y. Otsuka, and S. Yunoki, *Sci. Rep.* **2**, 992 (2012).